SISSA/20/2000/EP

New Physics behind the Standard Model's door?*

A. Masiero and O. Vives

SISSA, Via Beirut 2–4, 34013 Trieste, Italy and

INFN, sez. di Trieste, Trieste, Italy

Abstract

We review the main reasons pushing us beyond the SM and we argue in favor of new physics at the electroweak scale (hence experimentally accessible at present or near–future machines). We focus on the appealing possibility that such new physics is given by a supersymmetric (SUSY) extension of the SM. We discuss the minimal case, Constrained Minimal Supersymmetric SM, and more general (maybe more natural) cases where some of the drastic assumptions of the CMSSM are dropped. In particular, in these lectures we focus on CP violation and its relation to flavor physics in the SUSY context. CP constrains the low–energy SUSY extensions of the SM, but, at the same time, it provides new powerful tool for indirect SUSY searches.

* Lectures given by A. Masiero at the International School on Subnuclear Physics, 37th Course: "Basics and Highlights in Fundamental Physics", Erice, Italy, 29 August–7 September 1999

1 Introduction

The success of the standard model (SM) predictions is remarkably high and, indeed, to some extent, even beyond what we theorists would have expected. A common view before LEP started operating was that some new physics related to the electroweak symmetry breaking should show up when precisions at the percent level on some electroweak observable could be reached. As we know, on the contrary, even reaching sensitivities better than the percent has not given rise to any indication of departure from the SM predictions. All that can be summarized in a powerful statement about the "low–energy" limit of any kind of new physics beyond the SM: such new physics has to reproduce the SM with great accuracy when we consider its limit at energy scales of the order of the electroweak scale.

The fact that with the SM we have a knowledge of fundamental interactions up to energies of $\mathcal{O}(100)$ GeV should not be underestimated: it represents a tremendous and astonishing success of our gauge theory approach in particle physics and it is clear that it represents one of the great achievements in a century of great conquests in physics. Having said that, we are now confronting ourselves with an embarrassing question: if the SM is so extraordinarily good, does it make sense do go beyond it? The answer, in our view, is certainly positive. This "yes" is not only motivated by what we could define "philosophical" reasons, but there are specific motivations pushing us beyond the SM: we will group them in two broad categories, theoretical and "observational" reasons.

1.1 Theoretical reasons for new physics

The major theoretical conundra of the SM are related the following issues: flavor problem, unification of the fundamental interactions and gauge hierarchy problem. We briefly remind what they are about.

Flavor problem. All the masses and mixings of fermions are just free (unpredicted) parameters in the SM. To be sure, there is not even any hint in the SM about the number and rationale of fermion families. Leaving aside predictions for individual masses, not even a rough relation among fermion masses within the same generation or among different generations is present.

Unification of forces. At the time of the Fermi theory we had two couplings to describe the electromagnetic and the weak interactions (the electric constant and the Fermi constant, respectively). In the SM we are trading off those two couplings with two new couplings, the gauge couplings of $SU(2)$ and $U(1)$. Moreover, the gauge coupling of the strong interactions is very different from the other two. We cannot say that the SM represents a true unification of fundamental interactions, even leaving aside the problem that gravity is not considered at all by the model.

Gauge hierarchy. Fermion and vector boson masses are "protected" by symmetries in the SM (i.e., their masses can arise only when we break certain symmetries). On the contrary, the Higgs scalar mass does not enjoy such a symmetry protection. We would expect such mass to naturally jump to some higher scale where new physics sets in (this new energy scale could be some grand unification scale or the Planck mass, for instance). The only way to keep the Higgs mass at the electroweak scale is to perform incredibly accurate fine tunings of the parameters of the scalar sector.

1.2 "Observational" reasons for new physics

We have already said that all the experimental particle physics results of these last years have marked one success after the other of the SM. What do we mean then by "observational" difficulties for the SM? It is curious that such difficulties do not arise from observations within the strict particle physics domain, but rather they originate from possible **"clashes"** of the particle physics SM with the standard model of cosmology (i.e., the Hot Big Bang) or the standard model of the Sun. Explicitly we have in mind the following points.

Dark Matter. Denoting with Ω the ratio of the energy density to the critical energy density, the problem of dark matter (DM) can be summarized in the following two numbers: $\Omega_{DM} = 0.3$ and $\Omega_B < 0.1$. The first number denotes the amount of the contribution to Ω due to DM as inferred from measurements at the level of cluster of galaxies. The upper bound denotes the highest contribution of baryonic matter to Ω to have compatibility with one of the main pillars of the Big Bang model: nucleosynthesis. The clash between the above two numbers underlines the fact that we definitely need a large amount of non–baryonic DM. In the electroweak SM, no viable non–baryonic candidate exists to fulfill this task (remember that in the SM neutrinos are strictly massless). Hence, the existence of a (large) amount of non–baryonic DM pushes us to introduce new particles in addition to those of the SM.

Baryogenesis. Given that we have strong evidence that the Universe is vastly matter–antimatter asymmetric (i.e. no sizeable amount of primordial antimatter has survived), it is appealing to have a dynamical mechanism to give rise to such large baryon–antibaryon asymmetry starting from a symmetric situation. In the SM it is not possible to have such an efficient mechanism for baryogenesis. In spite of the fact that at the quantum level sphaleronic interactions violate baryon number in the SM, such violation cannot lead to the observed large matter–antimatter asymmetry (both CP violation is too tiny in the SM and, also, the present experimental lower bounds on the Higgs mass do not allow for a conveniently strong electroweak phase transition). Hence, a dynamical baryogenesis calls for the presence of new particles and interactions beyond the SM (successful mechanisms for baryogenesis in the context of new physics beyond the SM are well known).

Inflation. Several serious cosmological problems (flatness, causality, age of the Universe, ...) are beautifully solved if the early Universe underwent some period of exponential expansion (inflation). The SM with its Higgs doublet does not succeed to originate such an inflationary stage. Again some extensions of the SM, where in particular new scalar fields are introduced, are able to produce a temporary inflation of the early Universe.

1.3 The SM as an effective low–energy theory

The above theoretical and "observational" arguments strongly motivate us to go beyond the SM. On the other hand, the clear success of the SM in reproducing all the known phenomenology up to energies of the order of the electroweak scale is telling us that the SM has to be recovered as the low–energy limit of such new physics. Indeed, it may even well be the case that we have a "tower" of underlying theories which show up at different energy scales.

If we accept the above point of view, we may try to find signals of new physics considering the SM as a truncation to renormalizable operators of an effective low–energy theory which respects the $SU(3) \times SU(2) \times U(1)$ symmetry and whose fields are just those of the SM. The renormalizable (i.e. of canonical dimension less or equal to four) operators giving rise to the SM enjoy three crucial properties which have no reason to be shared by generic operators of dimension larger than four. They are the conservation (at any order in perturbation theory) of Baryon (B) and Lepton (L) numbers and an adequate suppression of Flavor Changing Neutral Current (FCNC) processes through the GIM mechanism.

Now consider the new physics (directly above the SM in the "tower" of new physics theories) to have a typical energy scale Λ. In the low–energy effective Lagrangian, such scale appears with a positive power only in the quadratic scalar term (scalar mass) and in the dimension zero operator which can be considered a cosmological constant. Notice that Λ cannot appear in dimension three operators related to fermion masses because chirality forbids direct fermion mass terms in the Lagrangian. Then, in all operators of dimension larger than four, Λ will show up in the denominator with powers increasing with the dimension of the corresponding operator.

The crucial question that all of us, theorists and experimentalists, ask ourselves is: where is Λ? Namely is it close to the electroweak scale (i.e. not much above 100 GeV) or is Λ of the order of the grand unification scale or the Planck scale? B– and L–violating processes and FCNC phenomena represent a potentially interesting clue to answer this fundamental question.

Take Λ to be close to the electroweak scale. Then we may expect non–renormalizable operators with B, L and flavor violations not to be largely suppressed by the presence of powers of Λ in the denominator. Actually this constitutes, in general, a formidable challenge for any model builder who wants to envisage new physics close to M_W. Theories with dynamical breaking of the electroweak symmetry (technicolour) and low–energy supersymmetry constitute examples of new physics with a "small" Λ. In these lectures we will see that the above general considerations on potentially large B, L and flavor violations apply to the SUSY case (it is well–known that FCNC represent a major problem also in technicolour schemes).

Alternatively, given the abovementioned potential danger of having a small Λ, one may feel it safer to send Λ to super–large values. Apart from kind of "philosophical" objections related to the unprecedented gap of many orders of magnitude without any new physics, the above discussion points out a typical problem of this approach. Since the quadratic scalar terms have a coefficient in front scaling with Λ^2, we expect all scalar masses to be of the order of the super–large scale Λ. This is the gauge hierarchy problem, and it constitutes the main (if not only) reason to believe that SUSY should be a low–energy symmetry.

Notice that the fact that SUSY should be a fundamental symmetry of Nature (something of which we have little doubt given the "beauty" of this symmetry) does not imply by any means that SUSY should be a low–energy symmetry, namely that it should hold unbroken down to the electroweak scale. SUSY may well be present in Nature but be broken at some very large scale (Planck scale or string compactification scale). In that case SUSY would be of no use in tackling the gauge hierarchy problem and its phenomenological relevance would be practically zero. On the other hand, if we invoke SUSY to tame the growth of the scalar mass terms with the scale Λ, then we are forced to take the view that SUSY should hold as a good symmetry down to a scale Λ close to the electroweak scale. Then B, L and FCNC may be useful for us to shed some light on the properties of the underlying theory from which the low–energy SUSY Lagrangian resulted. Let us add that there is an independent argument in favor of this view that SUSY should be a low–energy symmetry. The presence of SUSY partners at low energy creates the conditions to have a correct unification of the strong and electroweak interactions. If they were at M_{Planck} and the SM were all the physics up to super–large scales, the program of achieving such a unification would largely fail, unless one complicates the non–SUSY GUT scheme with a large number of Higgs representations and/or a breaking chain with intermediate mass scales is invoked.

In the above discussion, we stressed that we are not only insisting on the fact that SUSY should be present at some stage in Nature, but we are asking for something much more ambitious: we are asking for SUSY to be a low–energy symmetry, namely it should be broken at an energy scale as low as the electroweak symmetry breaking scale. This fact can never be overestimated. There are indeed several reasons pushing us to introduce SUSY: it is the most general symmetry compatible with a local, relativistic quantum field theory, it softens the degree of divergence of the theory, it looks promising for a consistent quantum description of gravity together with the other fundamental interactions. However, all these reasons are not telling us where we should expect SUSY to be broken. For that matter, we could even envisage the maybe “natural” possibility that SUSY is broken at the Planck scale. What is relevant for phenomenology is that the gauge hierarchy problem and, to some extent, the unification of the gauge couplings are actually forcing us to ask for SUSY to be unbroken down to the electroweak scale, hence implying that the SUSY copy of all the known particles, the so–called s–particles should have a mass in the 100–1000 GeV mass range. If LEP and Tevatron are not going to see any SUSY particle, at least the advent of LHC will be decisive in establishing whether low–energy SUSY actually exists or it is just a fruit of our (ingenious) speculations. Although even after LHC, in case of a negative result for the search of SUSY particles, we will not be able to “mathematically” exclude all the points of the SUSY parameter space, we will certainly be able to very reasonably assess whether the low–energy SUSY proposal makes sense or not.

Before the LHC (and maybe Tevatron) direct searches for SUSY signals we should ask ourselves whether we can hope to have some indirect manifestation of SUSY through virtual effects of the SUSY particles.

We know that, in the past, virtual effects (i.e. effects due to the exchange of yet unseen particles in the loops) were precious in leading us to major discoveries, like the prediction of the existence of the charm quark or the heaviness of the top quark long before its direct experimental observation. Here we focus on the potentialities of SUSY virtual effects in processes which are particluraly suppressed (or sometime even forbidden) in the SM; the flavor changing neutral current phenomena and the processes where CP violation is violated.

2 Flavor, CP and New Physics

The generation of fermion masses and mixings (“flavor problem”) gives rise to a first and important distinction among theories of new physics beyond the electroweak standard model.

One may conceive a kind of new physics which is completely “flavor blind”, i.e. new interactions which have nothing to do with the flavor structure. To provide an example of such a situation, consider a scheme where flavor arises at a very large scale (for instance the Planck mass) while new physics is represented by a supersymmetric extension of the SM with supersymmetry broken at a much lower scale and with the SUSY breaking transmitted to the observable sector by flavor–blind gauge interactions. In this case, one may think that the new physics does not cause any major change to the original flavor structure of the SM, namely that the pattern of fermion masses and mixings is compatible with the numerous and demanding tests of flavor changing neutral currents.

Alternatively, one can conceive a new physics which is entangled with the flavor problem. As an example consider a technicolour scheme where fermion masses and mixings arise through the exchange of new gauge bosons which mix together ordinary fermions and technifermions. Here we expect (correctly enough) new physics to have potential problems in accommodating the usual fermion spectrum with the adequate suppression of FCNC. As another example of new physics which is not flavor blind, take a more conventional SUSY model which is derived from a spontaneously broken N=1 supergravity and where the SUSY breaking information is conveyed to the ordinary sector of the theory through gravitational interactions. In this case we may expect that the scale at which flavor arises and the scale of SUSY breaking are not so different and possibly the mechanism itself of SUSY breaking and transmission is flavor–dependent. Under these circumstances, we may expect a potential flavor problem to arise, namely that SUSY contributions to FCNC processes are too large.

2.1 The Flavor Problem in SUSY

The potentiality of probing SUSY in FCNC phenomena was readily realized when the era of SUSY phenomenology started in the early 80's [1]. In particular, the major implication that the scalar partners of quarks of the same electric charge but belonging to different generations had to share a remarkably high mass degeneracy was emphasized.

Throughout the large amount of work in this last decade, it became clearer and clearer that generically talking of the implications of low–energy SUSY on FCNC may be rather misleading. Even in the Minimal SUSY extension of the SM (MSSM) [2] from the point of view of the particle content, we have a host of different situations. The so–called Constrained Minimal Supersymmetric Standard Model (CMSSM) is the simplest possibility, and the FCNC contributions can be computed in terms of a very limited set of unknown new SUSY parameters. Remarkably enough, this minimal model succeeds to pass all the set of FCNC tests unscathed. To be sure, it is possible to severely constrain the SUSY parameter space, for instance using $b \to s\gamma$, in a way which is complementary to what is achieved by direct SUSY searches at colliders.

However, the CMSSM is by no means equivalent to low–energy SUSY. A first sharp distinction concerns the mechanism of SUSY breaking and transmission to the observable sector which is chosen. As we mentioned above, in models with gauge–mediated SUSY breaking (GMSB models [3, 4, 5]) it may be possible to avoid the FCNC threat "ab initio" (notice that this is not an automatic feature of this class of models, but it depends on the specific choice of the sector which transmits the SUSY breaking information, the so–called messenger sector). The other more "canonical" class of SUSY theories (including also CMSSM) has gravitational messengers and a very large scale at which SUSY breaking occurs. In this talk we will focus only on this class of gravity–mediated SUSY breaking models. Even sticking to this more limited choice, we have a variety of options with very different implications for the flavor problem.

First, there exists an interesting large class of SUSY realizations where the customary R–parity (which is invoked to suppress proton decay) is replaced by other discrete symmetries which allow either baryon or lepton violating terms in the superpotential. But, even sticking to the more orthodox view of imposing R–parity, we are still left with a large variety of extensions of the MSSM at low energy. The point is that low–energy SUSY "feels" the new physics at the super–large scale at which supergravity (i.e., local supersymmetry) broke down. In this last couple of years, we have witnessed an increasing interest in supergravity realizations without the so–called flavor universality of the terms which break SUSY explicitly. Another class of low–energy SUSY realizations, which differ from the MSSM in the FCNC sector, is obtained from SUSY–GUT's. The interactions involving super–heavy particles in the energy range between the GUT and the Planck scale bear important implications for the amount and kind of FCNC that we expect at low energy.

Even when R–parity is imposed, the FCNC challenge is not over. It is true that in this case, analogously to what happens in the SM, no tree level FCNC contributions arise. However, it is well–known that this is a necessary but not sufficient condition to consider the FCNC problem overcome. The loop contributions to FCNC in the SM exhibit the presence of the GIM mechanism and we have to make sure that in the SUSY case with R parity some analog of the GIM mechanism is active.

To give a qualitative idea of what we mean by an effective super–GIM mechanism, let us consider the following simplified situation where the main features emerge clearly. Consider the SM box diagram responsible for the K^0–$\bar{K}^0$ mixing and take only two generations, i.e. only the up and charm quarks run in the loop. In this case, the GIM mechanism yields a suppression factor of $\mathcal{O}((m_c^2 - m_u^2)/M_W^2)$. If we replace the W boson and the up quarks in the loop with their SUSY partners and we take, for simplicity, all SUSY masses of the same order, we obtain a super–GIM factor which looks like the GIM one with the masses of the superparticles instead of those of the corresponding particles. The problem is that the up and charm squarks have masses which are much larger than those of the corresponding quarks. Hence the super–GIM factor tends to be of $\mathcal{O}(1)$ instead of being $\mathcal{O}(10^{-3})$ as it is in the SM case. To obtain this small number we would need a high degeneracy between the mass of the charm and up squarks. It is difficult to think that such a degeneracy may be accidental. After all, since we invoked SUSY for a naturalness problem (the gauge hierarchy issue), we should avoid invoking a fine–tuning to solve its problems! Then, one can turn to some symmetry reason. For instance, just sticking to this simple example that we are considering, one may think that the main bulk of the charm and up squark masses is the same, i.e. the mechanism of SUSY breaking should have some universality in providing the mass to these two squarks with the same electric charge. Flavor universality is by no means a prediction of low–energy SUSY. The absence of flavor universality of soft–breaking terms may result from radiative effects at the GUT scale or from effective supergravities derived from string theory. Indeed, from the point of view of these effective supergravity theories, it may appear more natural not to have such flavor universality. To obtain it one has to invoke particular circumstances, like, for instance, strong dilaton over moduli dominance in the breaking of supersymmetry, something which is certainly not expected on general ground.

Another possibility one may envisage is that the masses of the squarks are quite high, say above few TeV's. Then, even if they are not so degenerate in mass, the overall factor in front of the four–fermion operator responsible for the kaon mixing becomes smaller and smaller (it decreases quadratically with the mass of the squarks) and, consequently, one can respect the observational result. We see from this simple example that the issue of FCNC may be closely linked to the crucial problem of the way we break SUSY.

We now turn to some general remarks about the worries and hopes that CP violation arises in the SUSY context.

2.2 CP Violation in SUSY

CP violation has major potentialities to exhibit manifestations of new physics beyond the standard model. Indeed, it is quite a general feature that new physics possesses new CP violating phases in addition to the Cabibbo–Kobayashi–Maskawa (CKM) phase ($\delta_{\rm CKM}$) or, even in those cases where this does not occur, $\delta_{\rm CKM}$ shows up in interactions of the new particles, hence with potential departures from the SM expectations. Moreover, although the SM is able to account for the observed CP violation in the kaon system, we cannot say that we have tested so far the SM predictions for CP violation. The detection of CP violation in B physics will constitute a crucial test of the standard CKM picture within the SM. Again, on general grounds, we expect new physics to provide departures from the SM CKM scenario for CP violation in B physics. A final remark on reasons that make us optimistic in having new physics playing a major role in CP violation concerns the matter–antimatter asymmetry in the universe. Starting from a baryon–antibaryon symmetric universe, the SM is unable to account for the observed baryon asymmetry. The presence of new CP–violating contributions when one goes beyond the SM looks crucial to produce an efficient mechanism for the generation of a satisfactory ΔB asymmetry.

The above considerations apply well to the new physics represented by low–energy supersymmetric extensions of the SM. Indeed, as we will see below, supersymmetry introduces CP violating phases in addition to $\delta_{\rm CKM}$ and, even if one envisages particular situations where such extra–phases vanish, the phase $\delta_{\rm CKM}$ itself leads to new CP–violating contributions in processes where SUSY particles are exchanged. CP violation in B decays has all potentialities to exhibit departures from the SM CKM picture in low–energy SUSY extensions, although, as we will discuss, the detectability of such deviations strongly depends on the regions of the SUSY parameter space under consideration.

In any MSSM, at least two new "genuine" SUSY CP–violating phases are present. They originate from the SUSY parameters μ, M, A and B. The first of these parameters is the dimensionful coefficient of the $H_u H_d$ term of the superpotential. The remaining three parameters are present in the sector that softly breaks the N=1 global SUSY. M denotes the common value of the gaugino masses, A is the trilinear scalar coupling, while B denotes the bilinear scalar coupling. In our notation, all these three parameters are dimensionful. The simplest way to see which combinations of the phases of these four parameters are physical [6] is to notice that for vanishing values of μ, M, A and B the theory possesses two additional symmetries [7]. Indeed, letting B and μ vanish, a $U(1)$ Peccei–Quinn symmetry originates, which in particular rotates H_u and H_d. If M, A and B are set to zero, the Lagrangian acquires a continuous $U(1)$ R symmetry. Then we can consider μ, M, A and B as spurions which break the $U(1)_{PQ}$ and $U(1)_R$ symmetries. In this way, the question concerning the number and nature of the meaningful phases translates into the problem of finding the independent combinations of the four parameters which are invariant under $U(1)_{PQ}$ and $U(1)_R$ and determining their independent phases. There are three such independent combinations, but only two of their phases are independent. We use here the commonly adopted choice:

$$\varphi_A = \arg\left(A^* M\right), \qquad \varphi_B = \arg\left(B^* M\right). \tag{1}$$

where also $\arg\left(B\mu\right) = 0$, i.e. $\varphi_\mu = -\varphi_B$.

The main constraints on φ_A and φ_B come from their contribution to the electric dipole moments of the neutron and of the electron. For instance, the effect of φ_A and φ_B on the electric and chromoelectric dipole moments of the light quarks (u, d, s) lead to a contribution to d^e_N of order [8]

$$d^e_N \sim 2\left(\frac{100\mathrm{GeV}}{\tilde{m}}\right)^2 \sin\varphi_{A,B} \times 10^{-23} \mathrm{e\,cm}, \tag{2}$$

where $\tilde{m}$ here denotes a common mass for squarks and gluinos. The present experimental bound, $d^e_N < 1.1^{-25}$ e cm, implies that $\varphi_{A,B}$ should be $< 10^{-2}$, unless one pushes SUSY masses up to $\mathcal{O}(1$ TeV). A possible caveat to such an argument calling for a fine–tuning of $\varphi_{A,B}$ is that uncertainties in the estimate of the hadronic matrix elements could relax the severe bound in Eq. (2) [9].

In view of the previous considerations, most authors dealing with the MSSM prefer to simply put φ_A and φ_B equal to zero. Actually, one may argue in favor of this choice by considering the soft breaking sector of the MSSM as resulting from SUSY breaking mechanisms which force φ_A and φ_B to vanish. For instance, it is conceivable that both A and M originate from one same source of $U(1)_R$ breaking. Since φ_A "measures" the relative phase of A and M, in this case it would "naturally"vanish. In some specific models, it has been shown [10] that through an analogous mechanism also φ_B may vanish.

If $\varphi_A = \varphi_B = 0$, then the novelty of SUSY in CP violating contributions merely arises from the presence of the CKM phase in loops where SUSY particles run [11]. The crucial point is that the usual GIM suppression, which plays a major role in evaluating ε_K and ε'/ε in the SM, in the MSSM case (or more exactly in the CMSSM) is replaced by a super–GIM cancellation which has the same "power" of suppression as the original GIM (see previous section). Again, also in the CMSSM, as it is the case in the SM, the smallness of ε_K and ε'/ε is guaranteed not by the smallness of $\delta_{\rm CKM}$, but rather by the small CKM angles and/or small Yukawa couplings. By the same token, we do not expect any significant departure of the CMSSM from the SM predictions also concerning CP violation in B physics. As a matter of fact, given the large lower bounds on squark and gluino masses, one expects relatively tiny contributions of the SUSY loops in ε_K or ε'/ε in comparison with the normal W loops of the SM. Let us be more detailed on this point.

In the CMSSM, the gluino exchange contribution to FCNC is subleading with respect to chargino ($\chi^{\pm}$) and charged Higgs ($H^{\pm}$) exchanges. Hence, when dealing with CP violating FCNC processes in the CMSSM with $\varphi_A = \varphi_B = 0$, one can confine the analysis to $\chi^{\pm}$ and $H^{\pm}$ loops. If one takes all squarks to be degenerate in mass and heavier than ~ 200 GeV, then $\chi^{\pm} - \tilde{q}$ loops are obviously severely penalized with respect to the SM W^+–q loops (remember that at the vertices the same CKM angles occur in both cases).

The only chance for the CMSSM to produce some sizeable departure from the SM situation in CP violation is in the particular region of the parameter space where one has light $\tilde{q}$, $\chi^{\pm}$ and/or $H^{\pm}$. The best candidate (indeed the only one unless $\tan\beta \sim m_t/m_b$) for a light squark is the stop. Hence one can ask the following question: can the CMSSM present some novelties in CP–violating phenomena when we consider χ^+–$\tilde{t}$ loops with light $\tilde{t}$, χ^+ and/or H^+?

Several analyses in the literature tackle the above question or, to be more precise, the more general problem of the effect of light $\tilde{t}$ and χ^+ on FCNC processes [12, 13, 14]. A first important observation concerns the relative sign of the W^+–t loop with respect to the χ^+–$\tilde{t}$ and H^+–t contributions. As it is well known, the latter contribution always interferes positively with the SM one. Interestingly enough, in the region of the MSSM parameter space that we consider here, also the χ^+–$\tilde{t}$ contribution interferes constructively with the SM contribution. The second point regards the composition of the lightest chargino, i.e. whether the gaugino or higgsino component prevails. This is crucial since the light stop is predominantly $\tilde{t}_R$ and, hence, if the lightest chargino is mainly a wino, it couples to $\tilde{t}_R$ mostly through the LR mixing in the stop sector. Consequently, a suppression in the contribution to box diagrams going as $\sin^4\theta_{LR}$ is present (θ_{LR} denotes the mixing angle between the lighter and heavier stops). On the other hand, if the lightest chargino is predominantly a higgsino (i.e. $M_2 \gg \mu$ in the chargino mass matrix), then the χ^+–lighter $\tilde{t}$ contribution grows. In this case, contributions $\propto \theta_{LR}$ become negligible and, moreover, it can be shown that they are independent on the sign of μ. A detailed study is provided in reference [13, 14]. For instance, for $M_2/\mu = 10$, they find that the inclusion of the SUSY contribution to the box diagrams doubles the usual SM contribution for values of the lighter $\tilde{t}$ mass up to 100–120 GeV, using $\tan\beta = 1.8$, $M_{H^+} = 100$ TeV, $m_\chi = 90$ GeV and the mass of the heavier $\tilde{t}$ of 250 GeV. However, if m_χ is pushed up to 300 GeV, the χ^+–$\tilde{t}$ loop yields a contribution which is roughly 3 times less than in the case $m_\chi = 90$ GeV, hence leading to negligible departures from the SM expectation. In the cases where the SUSY contributions are sizeable, one obtains relevant restrictions on the ρ and η parameters of the CKM matrix by making a fit of the parameters A, ρ and η of the CKM matrix and of the total loop contribution to the experimental values of ε_K and ΔM_{B_d}. For instance, in the above–mentioned case in which the SUSY loop contribution equals the SM W^+–t loop, hence giving a total loop contribution which is twice as large as in the pure SM case, combining the ε_K and ΔM_{B_d} constraints leads to a region in the ρ–η plane with $0.15 < \rho < 0.40$ and $0.18 < \eta < 0.32$, excluding negative values of ρ.

In conclusion, the situation concerning CP violation in the MSSM case with $\varphi_A = \varphi_B = 0$ and exact universality in the soft–breaking sector can be summarized in the following way: the MSSM does not lead to any significant deviation from the SM expectation for CP–violating phenomena as d_N^e, ε_K, ε'/ε and CP violation in B physics; the only exception to this statement concerns a small portion of the MSSM parameter space where a very light $\tilde{t}$

($m_{\tilde{t}} < 100$ GeV) and χ^+ ($m_\chi \sim 90$ GeV) are present. In this latter particular situation, sizeable SUSY contributions to ε_K are possible and, consequently, major restrictions in the ρ–η plane can be inferred. Obviously, CP violation in B physics becomes a crucial test for this MSSM case with very light $\tilde{t}$ and χ^+. Interestingly enough, such low values of SUSY masses are at the border of the detectability region at LEP II.

In next Section, we will move to the case where, still keeping the minimality of the model, we switch on the new CP violating phases. Later on we will give up also the strict minimality related to the absence of new flavor structure in the SUSY breaking sector and we will see that, in those more general contexts, we can expect SUSY to significantly depart from the SM predictions in CP violating phenomena.

3 Flavor Blind SUSY Breaking and CP Violation

We have seen in the previous section that in any MSSM there are additional phases which can cause deviations from the predictions of the SM in CP violation experiments. In fact, in the CMSSM, there are already two new phases present, Eq.(1), and for most of the MSSM parameter space, the experimental bounds on the electric dipole moments (EDM) of the electron and neutron constrain these phases to be at most $\mathcal{O}(10^{-2})$. However, in the last few years, the possibility of having non–zero SUSY phases has again attracted a great deal of attention. Several new mechanisms have been proposed to suppress supersymmetric contributions to EDMs below the experimental bounds while allowing SUSY phases $\mathcal{O}(1)$. Methods of suppressing the EDMs consist of cancellation of various SUSY contributions among themselves [15], non universality of the soft breaking parameters at the unification scale [16] and approximately degenerate heavy sfermions for the first two generations [17]. In the presence of one of these mechanisms, large supersymmetric phases are naturally expected and EDMs should be generally close to the experimental bounds. [1]

In this section we will study the effects of these phases in CP violation

[1] In a more general (and maybe more natural) MSSM there are many other CP violating phases [18] that contribute to CP violating observables.

observables as ε_K, ε'/ε and B^0 CP asymmetries. In particular we will show that the presence of large susy phases is not enough to produce sizeable supersymmetric contributions to these observables. In fact, *in the absence of the CKM phase, a general MSSM with all possible phases in the soft–breaking terms, but no new flavor structure beyond the usual Yukawa matrices, can never give a sizeable contribution to ε_K, ε'/ε or hadronic B^0 CP asymmetries.* However, we will see in the next section, that as soon as one introduces some new flavor structure in the soft Susy–breaking sector, even if the CP violating phases are flavor independent, it is indeed possible to get sizeable CP contribution for large Susy phases and $\delta_{CKM} = 0$. Then, we can rephrase our sentence above in a different way: *A new result in hadronic B^0 CP asymmetries in the framework of supersymmetry would be a direct proof of the existence of a completely new flavor structure in the soft–breaking terms.* This means that B–factories will probe the flavor structure of the supersymmetry soft–breaking terms even before the direct discovery of the supersymmetric partners [19].

To prove this we will consider any MSSM, i.e. with the minimal supersymmetric particle content, with general **complex** soft–breaking terms, but with a flavor structure strictly given by the two familiar Yukawa matrices or any matrix strictly proportional to them. In these conditions, the most general structure of the soft–breaking terms at the large scale, that we call M_{GUT}, is,

$$\begin{array}{c}
(m_Q^2)_{ij} = m_Q^2\ \delta_{ij} \quad (m_U^2)_{ij} = m_U^2\ \delta_{ij} \quad (m_D^2)_{ij} = m_D^2\ \delta_{ij} \\
(m_L^2)_{ij} = m_L^2\ \delta_{ij} \quad (m_E^2)_{ij} = m_E^2\ \delta_{ij} \quad m_{H_1}^2 \quad m_{H_2}^2 \\
m_{\tilde{g}}\ e^{i\varphi_3} \quad m_{\tilde{W}}\ e^{i\varphi_2} \quad m_{\tilde{B}}\ e^{i\varphi_1} \quad (A_U)_{ij} = A_U\ e^{i\varphi_{A_U}}\ (Y_U)_{ij} \\
(A_D)_{ij} = A_D\ e^{i\varphi_{A_D}}\ (Y_D)_{ij} \quad (A_E)_{ij} = A_E\ e^{i\varphi_{A_E}}\ (Y_E)_{ij}.
\end{array} \tag{3}$$

where all the allowed phases are explicitly written and one of them can be removed by an R–rotation. All other numbers or matrices in this equation are always real. Notice that this structure covers, not only the CMSSM [20], but also most of Type I string motivated models considered so far from phenomenology [21, 22], gauge mediated models [3, 4, 5], minimal effective supersymmetry models [23, 24], etc.

Experiments of CP violation in the K or B systems only involve supersymmetric particles as virtual particles in the loops. This means that the phases in the soft–breaking terms can only appear in these experiments through the mass matrices of the SUSY particles. Then, the key point in our discussion will be the role played by the SUSY phases and the soft–breaking terms flavor structure in the low–energy sparticle mass matrices.

It is important to notice that, even in a model with flavor–universal soft–breaking terms at some high energy scale, as this is the case, some off–diagonality in the squark mass matrices appears at the electroweak scale. Working on the the so–called Super CKM basis (SCKM), where squarks are rotated parallel to the quarks so that Yukawa matrices are diagonalized, the squark mass matrix is not flavor diagonal at M_W. This is due to the fact that at M_{GUT} there are always two non–trivial flavor structures, namely the two Yukawa matrices for the up and down quarks, not simultaneously diagonalizable. This implies that through RGE evolution some flavor mixing leaks into the sfermion mass matrices. In a general Supersymmetric model, the presence of new flavor structures in the soft breaking terms would generate large flavor mixing in the sfermion mass matrices. However, in the CMSSM, the two Yukawa matrices are the only source of flavor change. Always in the SCKM basis, any off–diagonal entry in the sfermion mass matrices at M_W will be necessarily proportional to a product of Yukawa couplings. Then, a typical estimate for the element (i,j) in the L–L down squark mass matrix at the electroweak scale would necessarily be (see [20] for details),

$$(m^{2\,(D)}_{LL})_{ij} \approx \; c \; m_Q^2 \; Y^u_{ik} Y^{u\,*}_{jk}, \tag{4}$$

with c a proportionality factor between 0.1 and 1. This rough estimate provides the order of magnitude of the different entries in the sfermion mass matrices. It is important to notice that if the phases of these elements were $\mathcal{O}(1)$, due to some of the phases in equation (3), we would be able to give sizeable contributions, or even saturate, the different CP observables [25]. Then, it is clear that the relevant question for CP violation experiments is the presence of imaginary parts in these off–diagonal entries.

As explained in [20, 26, 27], once we have solved the Yukawa RGEs, the RGE equations of all soft–breaking terms are a set of linear differential equations. Then, they can be solved as a linear function of the initial conditions. For instance the scalar masses are,

$$\begin{aligned} m_S^2(M_W) = & \textstyle\sum_i \eta_S^{(\phi_i)} m_{\phi_i}^2 + \sum_{i\neq j} \left(\eta_S^{(g_i g_j)} e^{i(\varphi_i - \varphi_j)} + \eta_S^{(g_i g_j)\,T} e^{-i(\varphi_i-\varphi_j)} \right) m_{g_i} m_{g_j} \\ & + \textstyle\sum_i \eta_S^{(g_i)}\, m_{g_i}^2 + \sum_{ij} \left(\eta_S^{(g_i A_j)} e^{i(\varphi_i - \varphi_{A_j})} + \eta_S^{(g_i A_j)\,T} e^{-i(\varphi_i - \varphi_{A_j})} \right) m_{g_i} A_j \end{aligned}$$

$$+\sum_i \eta_S^{(A_i)} A_i^2 + \sum_{i\neq j} \left(\eta_S^{(A_i A_j)} e^{i(\varphi_{A_i}-\varphi_{A_j})} + \eta_S^{(A_i A_j)\,T} e^{-i(\varphi_{A_i}-\varphi_{A_j})}\right) A_i A_j \tag{5}$$

where $S = Q, U, D$, ϕ_i refers to any scalar, g_i to the different gauginos and A_i to any tri–linear coupling. In this equation, the different η matrices are 3×3 matrices, **strictly real** and all the allowed phases have been explicitly written. Regarding the imaginary parts, due to the hermiticity of the sfermion mass matrices, any imaginary part will always be associated to the non–symmetric part of the $\eta_S^{(g_i g_j)}$, $\eta_S^{(A_i A_j)}$ or $\eta_S^{(g_i A_j)}$ matrices. To estimate the size of these anti–symmetric parts, we can go to the RGE equations for the scalar mass matrices, where we use the same conventions and notation as in [20, 26]. Taking advantage of the linearity of these equations, we can directly write the evolution of the anti–symmetric parts, for instance $\hat{m}_Q^2 = m_Q^2 - (m_Q^2)^T$, as,

$$\begin{aligned}\frac{d\hat{m}_Q^2}{dt} = & -[\frac{1}{2}(\tilde{Y}_U\tilde{Y}_U^\dagger + \tilde{Y}_D\tilde{Y}_D^\dagger)\hat{m}_Q^2 + \frac{1}{2}\hat{m}_Q^2(\tilde{Y}_U\tilde{Y}_U^\dagger + \tilde{Y}_D\tilde{Y}_D^\dagger) + \\ & 2\ i\ \mathrm{Im}\{\tilde{A}_U\tilde{A}_U^\dagger + \tilde{A}_D\tilde{A}_D^\dagger\} + \tilde{Y}_U\hat{m}_U^2\tilde{Y}_U^\dagger + \tilde{Y}_D\hat{m}_D^2\tilde{Y}_D^\dagger]\end{aligned} \tag{6}$$

where, due to the reality of Yukawa matrices, we have used $Y^T = Y^\dagger$, and following [26] a tilde over the couplings ($\tilde{Y}$, $\tilde{A}$, ...) denotes a re–scaling by a factor $1/(4\pi)$. The evolution of the R–R squark mass matrices, m_U^2 and m_D^2, is completely analogous. With the initial conditions in equation (3), $\hat{m}_Q^2$, $\hat{m}_U^2$ and $\hat{m}_D^2$ at M_{GUT} are identically zero. Then, we can safely neglect the last two terms in equation (6) because they will only be a second order effect. This means that the only source for $\hat{m}_Q^2$ in equation (6) is necessarily $\mathrm{Im}\{A_U A_U^\dagger + A_D A_D^\dagger\}$ (also for $\hat{m}_{U,D}^2$) .

The next step is then to analyze the RGE for the tri–linear couplings,

$$\begin{aligned}\frac{d\tilde{A}_U}{dt} = & \frac{1}{2}\Big(\frac{16}{3}\tilde{\alpha}_3 + 3\tilde{\alpha}_2 + \frac{1}{9}\tilde{\alpha}_1\Big)\tilde{A}_U - \Big(\frac{16}{3}\tilde{\alpha}_3 M_3 + 3\tilde{\alpha}_2 M_2 + \frac{1}{9}\tilde{\alpha}_1 M_1\Big)\tilde{Y}_U - \\ & \Big(2\tilde{A}_U\tilde{Y}_U^\dagger\tilde{Y}_U + 3Tr(\tilde{A}_U\tilde{Y}_U^\dagger)\tilde{Y}_U + \frac{5}{2}\tilde{Y}_U\tilde{Y}_U^\dagger\tilde{A}_U + \frac{3}{2}Tr(\tilde{Y}_U\tilde{Y}_U^\dagger)\tilde{A}_U + \\ & \tilde{A}_D\tilde{Y}_D^\dagger\tilde{Y}_U + \frac{1}{2}\tilde{Y}_D\tilde{Y}_D^\dagger\tilde{A}_U\Big)\end{aligned} \tag{7}$$

with an equivalent equation for A_D. With the general initial conditions in equation (3), A_U is complex at any scale. However, we are interested in the imaginary parts of $A_U A_U^\dagger$. At M_{GUT} this combination is exactly real, but, due to different renormalization of different elements of the matrix, this is not true anymore at a different scale. Nevertheless, a careful analysis of equation (7) is enough to convince ourselves that these imaginary parts are extremely small. Let us, for a moment, neglect the terms involving $\tilde{A}_D \tilde{Y}_D^\dagger$ or $\tilde{Y}_D \tilde{Y}_D^\dagger$ from the above equation or, strictly speaking, from the complete set of MSSM RGE. Then, the only flavor structure appearing in equation (7) at M_{GUT} is Y_U. We can always go to the basis where Y_U is diagonal and then we will have A_U exactly diagonal at any scale. In particular this means that $\mathrm{Im}\{A_U A_U^\dagger\}$ would always exactly vanish. A completely parallel reasoning can be applied to A_D and $\mathrm{Im}\{A_D A_D^\dagger\}$. Hence, simply taking into account the flavor structure, our conclusion is that, necessarily, any non–vanishing element of $\mathrm{Im}[A_U A_U^\dagger + A_D A_D^\dagger]$ and hence of $\hat{m}_Q^2$ must be proportional to $(\tilde{Y}_D \tilde{Y}_D^\dagger \tilde{Y}_U \tilde{Y}_U^\dagger - H.C.)$. So, we can expect them to be,

$$\begin{aligned}
(\hat{m}_Q^2)_{i<j} &\approx K \left(Y_D Y_D^\dagger Y_U Y_U^\dagger - H.C. \right)_{i<j} \\
(\hat{m}_Q^2)_{12} &\approx K \cos^{-2}\beta \; (h_s h_t \lambda^5) \\
(\hat{m}_Q^2)_{13} &\approx K \cos^{-2}\beta \; (h_b h_t \lambda^3) \\
(\hat{m}_Q^2)_{23} &\approx K \cos^{-2}\beta \; (h_b h_t \lambda^2),
\end{aligned} \tag{8}$$

where $h_i = m_i^2/v^2$, with $v = \sqrt{v_1^2 + v_2^2}$ the vacuum expectation value of the Higgs, $\lambda = \sin\theta_c$ and K is a proportionality constant that includes the effects of the running from M_{GUT} to M_W. To estimate this constant, we have to keep in mind that the imaginary parts of $A_U A_U^\dagger$ are generated through the RGE running and then, these imaginary parts generate $\hat{m}_Q^2$ as a second order effect. This means that roughly $K \simeq \mathcal{O}(10^{-2})$ times a combination of initial conditions as in equation (5). So, we estimate these matrix elements to be $(\cos^{-2}\beta\{10^{-12}, 6\times 10^{-8}, 3\times 10^{-7}\})$ times initial conditions. This was exactly the result we found for the A–g terms in [20] in the framework of the CMSSM. In fact, now it is clear that this is the same for all the terms in equation (5), g_i–A_j, g_i–g_j and A_i–A_j, irrespectively of the presence of an arbitrary number of new phases. This discussion can be directly applied for the R–R matrices.

Hence, so far, we have shown that the L–L or R–R squark mass matrices are still essentially real. The only complex matrices, then, will still be the L–R matrices that include, from the very beginning, the phases φ_{A_i} and φ_μ. Once more, the size of these entries is determined by the Yukawa elements with these two phases providing the complex structure. In fact, we can follow the same reasoning used after Eq.(7). In the absence of the terms involving $\tilde{A}_D \tilde{Y}_D^\dagger$ or $\tilde{Y}_D \tilde{Y}_D^\dagger$, the $\tilde{A}_U$ matrix would be exactly diagonalized when we diagonalize the Yukawa matrices. So, any off–diagonal element in the SCKM basis will be proportional to three Yukawas, $(\tilde{Y}_D \tilde{Y}_D^\dagger \tilde{Y}_U)$ and hence sufficiently small. Notice that this situation is not new for these more general MSSM models and it was already present even in the CMSSM. We can conclude, then, that the structure of the sfermion mass matrices at M_W is not modified from the familiar structure already present in the CMSSM, irrespective of the presence of an arbitrary number of new SUSY phases.

3.1 Indirect CP violation

Next, we will consider indirect CP violation both in the K and B systems. In the SM, neutral meson mixing arises at one loop through the well–known W–box. However, in the MSSM, there are new contributions to $\Delta F = 2$ processes coming from boxes mediated by supersymmetric particles. These are: charged Higgs boxes ($H^\pm$), chargino boxes ($\chi^\pm$) and gluino–neutralino boxes ($\tilde{g}$, χ^0). $\mathcal{M}$–$\bar{\mathcal{M}}$ mixing is correctly described by the $\Delta F = 2$ effective Hamiltonian, $\mathcal{H}_{eff}^{\Delta F=2}$, which can be decomposed as,

$$\mathcal{H}_{eff}^{\Delta F=2} = -\frac{G_F^2 M_W^2}{(2\pi)^2} (K_{td}^* K_{tq})^2 (C_1(\mu) Q_1(\mu) + C_2(\mu) Q_2(\mu) + C_3(\mu) Q_3(\mu)). \tag{9}$$

With the relevant four–fermion operators given by

$$Q_1 = \bar{d}_L^\alpha \gamma^\mu q_L^\alpha \cdot \bar{d}_L^\beta \gamma_\mu q_L^\beta, \quad Q_2 = \bar{d}_L^\alpha q_R^\alpha \cdot \bar{d}_L^\beta q_R^\beta, \quad Q_3 = \bar{d}_L^\alpha q_R^\beta \cdot \bar{d}_L^\beta q_R^\alpha, \tag{10}$$

where $q = s, b$ for the K and B–systems respectively, α, β are color indices and K_{ij} is the CKM mixing matrix. In the CMSSM, these are the only three operators present in the limit of vanishing m_d. The Wilson coefficients, $C_1(\mu)$, $C_2(\mu)$ and $C_3(\mu)$, receive contributions from the different supersymmetric boxes,

$$\begin{aligned} C_1(M_W) &= C_1^W(M_W) + C_1^H(M_W) + C_1^{\tilde{g},\chi^0}(M_W) + C_1^\chi(M_W) \\ C_2(M_W) &= C_2^H(M_W) + C_2^{\tilde{g}}(M_W) \\ C_3(M_W) &= C_3^{\tilde{g},\chi^0}(M_W) + C_3^\chi(M_W) \end{aligned} \tag{11}$$

Both, the usual SM W–box and the charged Higgs box contribute to these operators. However, with $\delta_{CKM} = 0$, these contributions do not contain any complex phase and hence cannot generate an imaginary part for these Wilson coefficients.

Then, Gluino and neutralino contributions are specifically supersymmetric. They involve the superpartners of quarks and gauge bosons. Here, the source of flavor mixing is not directly the usual CKM matrix. It is the presence of off–diagonal elements in the sfermion mass matrices, as discussed in section 3. To analyze these contributions, it is convenient to use the so–called Mass Insertion (MI) approximation [28, 25]. To define the MI we go to the SCKM basis. In this basis, off–diagonal flavor–changing effects can be estimated by insertion of flavor–off–diagonal components of the mass–squared matrices. By normalizing by an average squark mass–squared $m^2_{\tilde{q}}$, we define,

$$(\delta^d_{LL})_{ij} = \frac{({m^2}^{(d)}_{LL})_{ij}}{m^2_{\tilde{q}}} \quad (\delta^d_{RR})_{ij} = \frac{({m^2}^{(d)}_{RR})_{ij}}{m^2_{\tilde{q}}} \quad (\delta^d_{LR})_{ij} = \frac{({m^2}^{(d)}_{LR})_{ij}}{m^2_{\tilde{q}}} \tag{12}$$

with ${m^2}^{(d)}_{AB}$ the squark mass matrices in the SCKM basis.

From the point of view of CP violation, we will always need a complex Wilson coefficient. In the SCKM basis all gluino vertices are flavor diagonal and real. Then, a complex MI in one of the sfermion lines is always required. Only L–L mass insertions enter at first order in the Wilson coefficient ${C_1}^{\tilde{g},\chi^0}(M_W)$. From equation (8), the imaginary parts of these MI are at most $\mathcal{O}(10^{-6})$ for the b–s transitions and smaller otherwise [20]. Comparing these values with the phenomenological bounds required to saturate the measured values of these processes [25] we can easily see that we are always several orders of magnitude below.

In the case of the Wilson coefficients ${C_2}^{\tilde{g}}(M_W)$ and ${C_3}^{\tilde{g}}(M_W)$, the involved MI are L–R. However, as explained in section 3, these MI are always suppressed by light masses of right handed squark plus two additional up Yukawas. Moreover, in the case of b–s transitions they are directly constrained by the $b \to s\gamma$ decay. Hence, gluino boxes, in the absence of new flavor structures, can never give sizeable contributions to indirect CP violation processes [20].

The chargino contributions to these Wilson coefficients were discussed in great detail in the CMSSM framework in reference [20]. In this more general MSSM, we find very similar results due to the absence of new flavor structure.

Basically, in the chargino boxes, flavor mixing comes explicitly from the CKM mixing matrix, although off–diagonality in the sfermion mass matrix introduces a small additional source of flavor mixing.

$$C_1^{\chi}(M_W) = \sum_{i,j=1}^{2} \sum_{k,l=1}^{6} \sum_{\alpha\gamma\alpha'\gamma'} \frac{K^*_{\alpha' d} K_{\alpha q} K^*_{\gamma' d} K_{\gamma q}}{(K^*_{td} K_{tq})^2}$$
$$[G^{(\alpha,k)i} G^{(\alpha',k)j*} G^{(\gamma',l)i*} G^{(\gamma,l)j}\ Y_1(z_k, z_l, s_i, s_j)] \qquad (13)$$

where $K_{\alpha q} G^{(\alpha,k)i}$ represent the coupling of chargino and squark k to left–handed down quark q, $z_k = M^2_{\tilde{u}_k}/M^2_W$ and $s_i = M^2_{\tilde{\chi}_i}/M^2_W$. The explicit expressions for the loop functions can be found in reference [20]. These couplings, in terms of the standard mixing matrices [29, 26],

$$G^{(\alpha,k)i} \quad = \quad \left(\Gamma^{k\alpha}_{UL} V^*_{i1} - \frac{m_\alpha}{\sqrt{2} M_W \sin\beta} \Gamma^{k\alpha}_{UR} V^*_{i2} \right). \qquad (14)$$

$G^{(\alpha,k)i}$ are in general complex, as both φ_μ and φ_{A_i} are present in the different mixing matrices.

The main part of C_1^{χ} in equation (13) will be given by pure CKM flavor mixing, neglecting the additional flavor mixing in the squark mass matrix [30, 14]. This means, $\alpha = \alpha'$ and $\gamma = \gamma'$. In these conditions, using the symmetry of the loop function $Y_1(a, b, c, d)$ under the exchange of any two indices it is easy to prove that C_1^{χ} would be exactly real [23]. This is not exactly true either in the CMSSM or in our more general MSSM, where there is additional flavor change in the sfermion mass matrices. Here, some imaginary parts appear in the C_1^{χ} in equation (13). In figure 1 we show in a scatter plot the size of imaginary and real parts of C_1^{χ} in the B system for a fixed value of $\tan\beta = 40$. We see that this Wilson coefficient is always real up to a part in 10^3. In any case, this is out of reach for the foreseen B–factories. For the K system, imaginary parts are still smaller due to smaller mixing angles with the stop.

Finally, chargino boxes contribute also to the quirality changing Wilson coefficient $C_3^{\chi}(M_W)$,

$$C_3^{\chi}(M_W) = \sum_{i,j=1}^{2} \sum_{k,l=1}^{6} \sum_{\alpha\gamma\alpha'\gamma'} \frac{K^*_{\alpha' d} K_{\alpha q} K^*_{\gamma' d} K_{\gamma q}}{(K^*_{td} K_{tq})^2} \ \frac{m_q^2}{2 M_W^2 \cos^2\beta}$$
$$H^{(\alpha,k)i} G^{(\alpha',k)j*} G^{(\gamma',l)i*} H^{(\gamma,l)j} Y_2(z_k, z_l, s_i, s_j) \qquad (15)$$

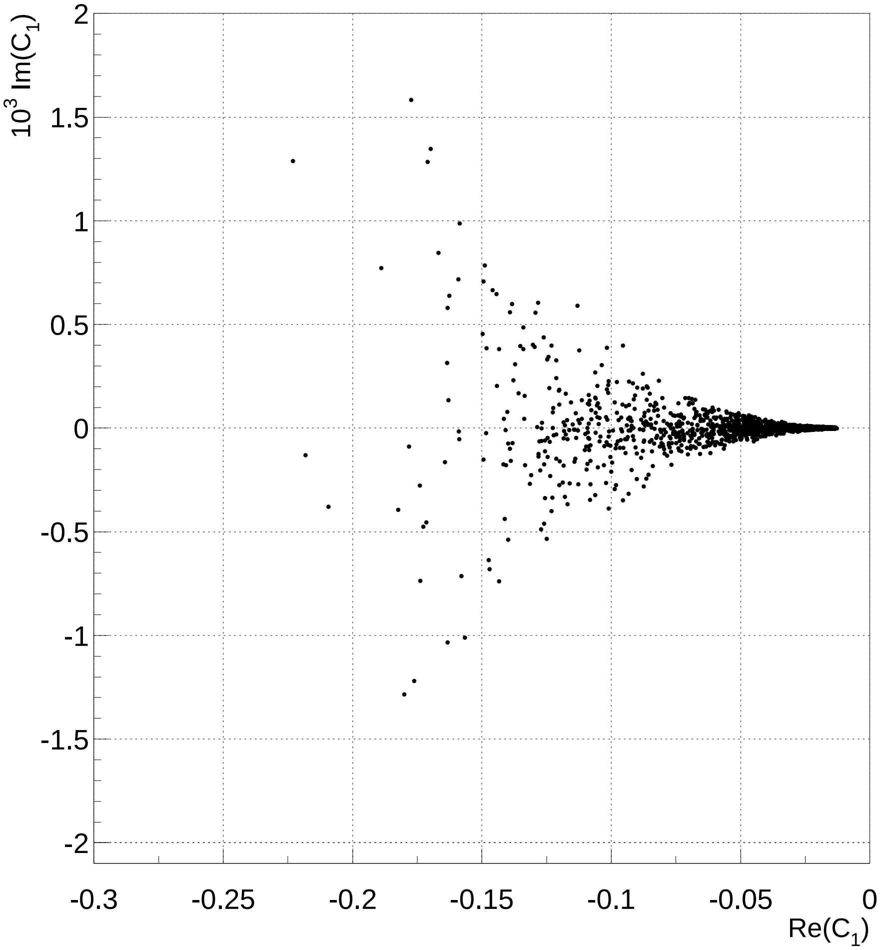

Figure 1: Imaginary and Real parts of the Wilson coefficient C_1^χ in B mixing.

where $m_q/(\sqrt{2}M_W\cos\beta)\cdot K_{\alpha q}\cdot H^{(\alpha,k)i}$ is the coupling of chargino and squark to the right–handed down quark q [29, 26],

$$H^{(\alpha,k)i} = -U_{i2}\Gamma_{UL}^{k\alpha}. \tag{16}$$

Unlike the C_1^{χ} Wilson coefficient, due to the differences between H and G couplings, C_3^{χ} is complex even in the absence of intergenerational mixing in the sfermion mass matrices [23]. Then, the presence of these small flavor violating entries in the up–squark mass matrix hardly modifies the results obtained in their absence [30, 14, 20]. In fact, in spite the presence of the Yukawa coupling squared, $m_q^2/(2M_W^2\cos^2\beta)$, this contribution could be relevant in the large $\tan\beta$ regime. For instance, in B^0–$\bar{B}^0$ mixing we have $m_b^2/(2M_W^2\cos^2\beta)$ that for $\tan\beta \gtrsim 25$ is larger than 1 and so, it is not suppressed at all when compared with the C_1^{χ} Wilson Coefficient. This means that this contribution can be very important in the large $\tan\beta$ regime [23] and could have observable effects in CP violation experiments in the new B–factories. However, we will show next that when we include the constraints coming from $b \to s\gamma$ these chargino contributions are also reduced to an unobservable level.

The chargino contributes to the $b \to s\gamma$ decay through the Wilson coefficients $\mathcal{C}_7$ and $\mathcal{C}_8$, corresponding to the photon and gluon dipole penguins respectively [26, 31, 20]. In the large $\tan\beta$ regime, we can approximate these Wilson coefficients as [20],

$$\begin{aligned}
\mathcal{C}_7^{\chi^\pm}(M_W) &= \textstyle\sum_{k=1}^{6}\sum_{i=1}^{2}\sum_{\alpha,\beta=u,c,t}\ \dfrac{K_{\alpha b}K_{\beta s}^*}{K_{tb}K_{ts}^*}\dfrac{m_b}{\sqrt{2}M_W\cos\beta} \\
&\qquad H^{(\alpha,k)i}G^{*(\beta,k)i}\ \frac{M_{\chi^i}}{m_b}\ F_R^7(z_k,s_i) \\
\mathcal{C}_8^{\chi^\pm}(M_W) &= \textstyle\sum_{k=1}^{6}\sum_{i=1}^{2}\sum_{\alpha,\beta=u,c,t}\ \dfrac{K_{\alpha b}K_{\beta s}^*}{K_{tb}K_{ts}^*}\dfrac{m_b}{\sqrt{2}M_W\cos\beta} \\
&\qquad H^{(\alpha,k)i}G^{*(\beta,k)i}\ \frac{M_{\chi^i}}{m_b}\ F_R^8(z_k,s_i)
\end{aligned} \tag{17}$$

Now, if we compare the chargino contributions to these Wilson coefficients and to the coefficient C_3, equations (15) and (17), we can see that they are deeply related. In fact, in the approximation where the two different loop functions involved are of the same order, we have,

$$C_3(M_W) \approx (\mathcal{C}_7(M_W))^2\frac{m_q^2}{M_W^2} \tag{18}$$

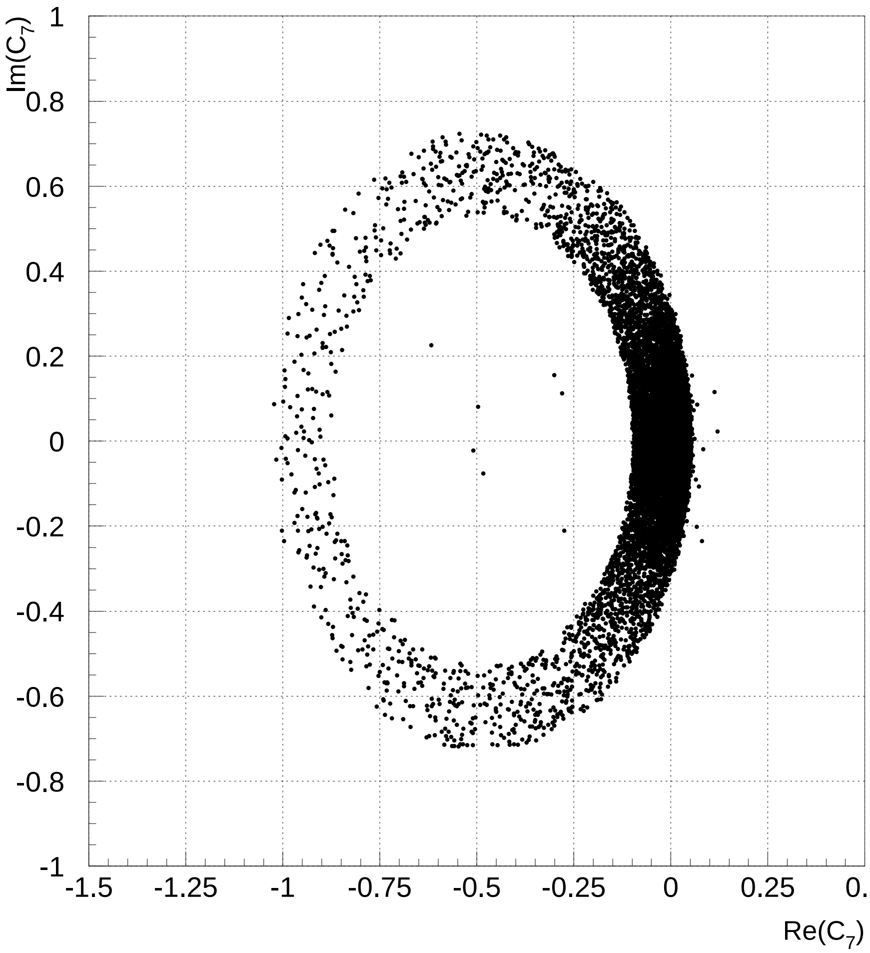

Figure 2: Experimental constraints on the Wilson Coefficient $\mathcal{C}_7$

In figure 2, we show a scatter plot of the allowed values of $Re(\mathcal{C}_7)$ versus $Im(\mathcal{C}_7)$ in the CMSSM for a fixed value of $\tan\beta = 40$ [20] with the constraints from the decay $B \to X_s\gamma$ taken from the reference [32]. Notice that a relatively large value of $\tan\beta$, for example, $\tan\beta \gtrsim 10$, is needed to compensate the W and charged Higgs contributions and cover the whole allowed area with positive and negative values. However, the shape of the plot is clearly independent of $\tan\beta$, only the number of allowed points and its location in the allowed area depend on the value considered. Then, figure 3 shows the allowed values for a re–scaled Wilson coefficient $\bar{C}_3(M_W) = M_W^2/m_q^2 C_3(M_W)$ corresponding to the same allowed points of the SUSY parameter space in figure 2. As we anticipated previously, the allowed values for $\bar{C}_3$ are close to the square of the values of $\mathcal{C}_7$ in figure 2 slightly scaled by different values of the loop functions.

We can immediately translate this result to a constraint on the size of the chargino contributions to $\varepsilon_{\mathcal{M}}$.

$$\begin{aligned} \varepsilon_{\mathcal{M}} \;=\; & \frac{G_F^2 M_W^2}{4\pi^2\sqrt{2}\;\Delta M_{\mathcal{M}}}\;\frac{(K_{td}K_{tq})^2}{24} F_{\mathcal{M}}^2\; M_{\mathcal{M}} \\ & \frac{M_{\mathcal{M}}^2}{m_q^2(\mu)+m_d^2(\mu)}\;\eta_3(\mu)\; B_3(\mu)\; Im[C_3] \end{aligned} \tag{19}$$

In this expression $M_{\mathcal{M}}$, $\Delta M_{\mathcal{M}}$ and $F_{\mathcal{M}}$ denote the mass, mass difference and decay constant of the neutral meson $\mathcal{M}^0$. The coefficient $\eta_3(\mu) = 2.93$ [33] includes the RGE effects from M_W to the meson mass scale, μ, and $B_3(\mu)$ is the B–parameter associated with the matrix element of the Q_3 operator [33].

For the K system, using the experimentally measured value of ΔM_K we obtain,

$$\varepsilon_K^{\chi} \;=\; 1.7\times 10^{-2}\;\frac{m_s^2}{M_W^2}\; Im[\bar{C}_3] \approx 0.4\times 10^{-7}\; Im[\bar{C}_3] \tag{20}$$

Given the allowed values of $\bar{C}_3$ in figure 3, this means that in the MSSM, even with large SUSY phases, chargino cannot produce a sizeable contribution to ε_K.

The case of B^0–$\bar{B}^0$ mixing has a particular interest due to the arrival of new data from the B–factories. In fact, in the large $\tan\beta$ regime chargino contributions to indirect CP violation can be very important. However, for

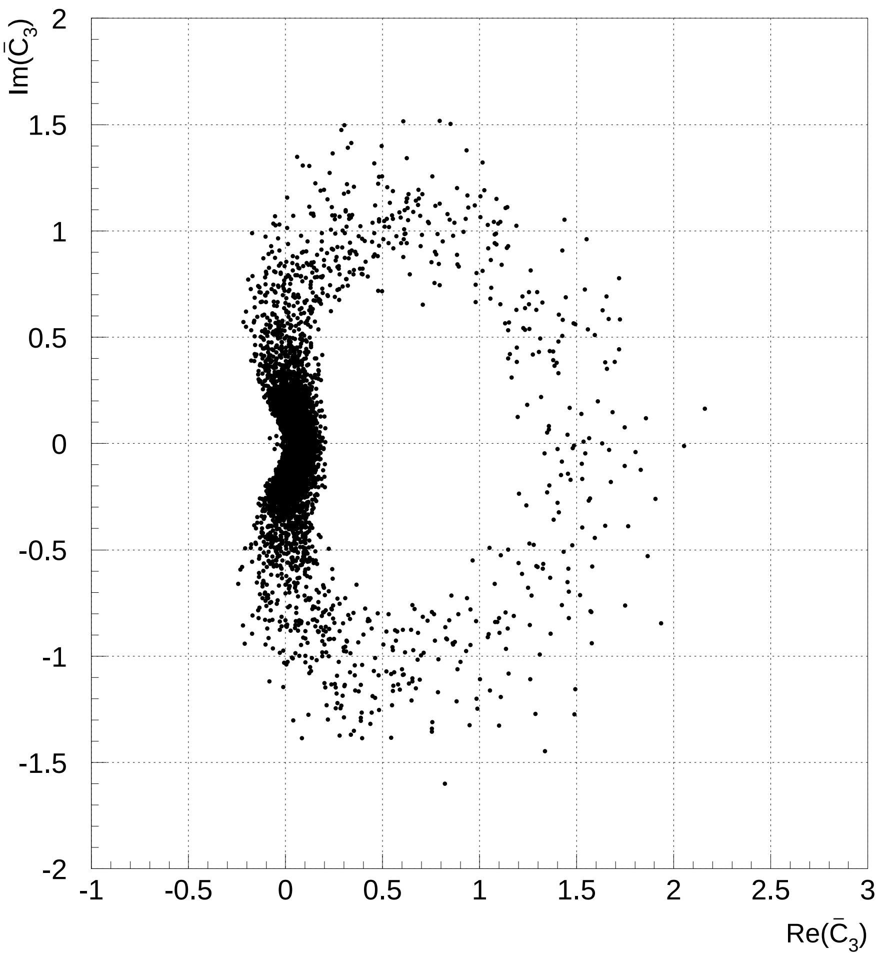

Figure 3: Allowed values for the re–scaled WC $\bar{C}_3$

any value of $\tan\beta$, we must satisfy the bounds from the $b \to s\gamma$ decay. Then, if we apply these constraints to the B^0–$\bar{B}^0$ mixing,

$$\varepsilon_B^{\chi} \; = \; 0.17 \, \frac{m_b^2}{M_W^2} \; Im[\bar{C}_3] \approx 0.5 \times 10^{-3} \; Im[\bar{C}_3] \tag{21}$$

where once again, with the allowed values of figure 3 we get a very small contribution to CP violation in the mixing. We must take into account that the mixing–induced CP phase, θ_M, measurable in B^0 CP asymmetries, is related to ε_B by $\theta_M = \arcsin\{2\sqrt{2} \cdot \varepsilon_B\}$. The expected sensitivities on the CP phases at the B factories are around ± 0.1 radians, so this supersymmetric chargino contribution will be absolutely out of reach.

3.2 Direct CP violation

To complete our analysis, we consider now direct CP violation. In this case, the different decay processes are described by a $\Delta F = 1$ effective Hamiltonian. A complete operator basis for these transitions in a general MSSM involves 14 different operators [25]. The main difference with the case of indirect CP violation is that these operators receive contributions both from box and penguin diagrams. Nevertheless, the discussion of the presence of imaginary parts is completely analogous to the case of indirect CP violation.

Once more, in the gluino case, L–L transitions are real to a very good approximation, and several orders of magnitude below the phenomenological bounds [25]. On the other hand, L–R transitions are suppressed by two up Yukawas and a down quark mass and $b \to s\gamma$ decay. This is always true for the squark mass matrices obtained in section 3, and valid both for boxes and penguins.

Finally, we are left with chargino contributions. The analysis of chargino boxes is exactly the same as in the previous section. In fact, even the Wilson coefficients are identical except some CKM elements that can always be factored out. Then, for the penguins, L–L transitions are exactly real if we neglect inter–generational mixing in the squark mass matrices. Taking into account this small mixing we find, for the very same reasons as in the indirect CP violation case, that imaginary parts are far too small. The relation of the $b \to s\gamma$ decay with the L–R chargino penguins is in this case even more transparent than for the boxes. So, our conclusion is again that no new supersymmetric CP violation effects are possible in ε'/ε or hadronic B CP asymmetries.

However, there is still one possibility to observe the effects of the new supersymmetric phases even in the absence of new flavor structure. We have seen that the reason for the smallness of the contributions of chargino L–R transitions is the experimental bound from the $B \to X_s\gamma$ branching ratio. This bound makes the chirality changing transitions, although complex, too small to compete with L–L transitions. Hence, in these conditions, just the processes where only chirality changing operators contribute (EDMs or $b \to s\gamma$), or observables where chirality flip operators are relevant ($b \to sl^+l^-$) can show the effects of new supersymmetric phases [24, 20].

4 CP Violation in the presence of new Flavor Structures

In section 3, we have shown that CP violation effects are always small in models with flavor blind soft–breaking terms. However, as soon as one introduces some new flavor structure in the soft breaking sector, it is indeed possible to get sizeable CP contribution for large Susy phases and $\delta_{CKM} = 0$ [16, 34, 35]. To show this, we will mainly concentrate in new supersymmetric contributions to ε'/ε.

In the CMSSM, the SUSY contribution to ε'/ε is small [36, 19]. However in a MSSM with a more general framework of flavor structure it is relatively easy to obtain larger SUSY effects to ε'/ε. In ref. [37] it was shown that such large SUSY contributions arise once one assumes that: i) hierarchical quark Yukawa matrices are protected by flavor symmetry, ii) a generic dependence of Yukawa matrices on Polonyi/moduli fields is present (as expected in many supergravity/superstring theories), iii) the Cabibbo rotation originates from the down–sector and iv) the phases are of order unity. In fact, in [37], it was illustrated how the observed ε'/ε could be mostly or entirely due to the SUSY contribution.

The universality of the breaking is a strong assumption and is known not to be true in many supergravity and string inspired models [38]. In these models, we expect at least some non–universality in the squark mass matrices or tri–linear terms at the supersymmetry breaking scale. Hence, sizeable flavor–off-diagonal entries will appear in the squark mass matrices. In this regard, gluino contributions to ε'/ε are especially sensitive to $(\delta^d_{12})_{LR}$; even $|\mathrm{Im}(\delta^d_{12})^2_{LR}| \sim 10^{-5}$ gives a significant contribution to ε'/ε while keeping the contributions from this MI to Δm_K and ε_K well bellow the phenomenological bounds. The situation is the opposite for L–L and R–R mass insertions; the stringent bounds on $(\delta^d_{12})_{LL}$ and $(\delta^d_{12})_{RR}$ from Δm_K and ε_K prevent them to contribute significantly to ε'/ε.

The LR squark mass matrix has the same flavor structure as the fermion Yukawa matrix and both, in fact, originate from the superpotential couplings. It may be appealing to invoke the presence of an underlying flavor symmetry restricting the form of the Yukawa matrices to explain their hierarchical forms. Then, the LR mass matrix is expected to have a very similar form as the Yukawa matrix. Indeed, we expect the components of the LR mass matrix to be roughly the SUSY breaking scale (e.g., the gravitino mass) times the corresponding component of the quark mass matrix. However, there is no reason for them to be simultaneously diagonalizable based on this general argument. To make an order of magnitude estimate, we take the down quark mass matrix for the first and second generations to be (following our assumption iii)),

$$Y^d v_1 \simeq \begin{pmatrix} m_d & m_s V_{us} \\ & m_s \end{pmatrix}, \tag{22}$$

where the (2,1) element is unknown due to our lack of knowledge on the mixings among right–handed quarks (if we neglect small terms $m_d V_{cd}$). Based on the general considerations on the LR mass matrix above, we expect

$$m^{2(d)}_{LR} \simeq m_{3/2} \begin{pmatrix} a m_d & b m_s V_{us} \\ & c m_s \end{pmatrix}, \tag{23}$$

where a, b, c are constants of order unity. Unless $a = b = c$ exactly, M_d and $m^{2,d}_{LR}$ are not simultaneously diagonalizable and we find

$$(\delta^d_{12})_{LR} \simeq \frac{m_{3/2} m_s V_{us}}{m^2_{\tilde{q}}} = 2 \times 10^{-5} \left(\frac{m_s(M_{Pl})}{50\ \mathrm{MeV}}\right) \left(\frac{m_{3/2}}{m_{\tilde{q}}}\right) \left(\frac{500\ \mathrm{GeV}}{\mathrm{m}_{\tilde{\mathrm{q}}}}\right). \tag{24}$$

It turns out that, following the simplest implementation along the lines of the above described idea, the amount of flavor changing LR mass insertion in the s and d–squark propagator results to roughly saturate the bound from ε'/ε if a SUSY phase of order unity is present [37].

This line of work has received a great deal of attention in recent times, after the last experimental measurements of ε'/ε in KTeV and NA31 [39, 40]. The effects of non–universal A terms in CP violation experiments were previously analyzed by Abel and Frere [41] and after this new measurement discussed in many different works [16]. In the following we show a complete realization of the above Masiero–Murayama (MM) mechanism from a Type I string–derived model recently presented by one of the authors [42].

4.1 Type I string model and ε'/ε

In first place we explain our starting model, which is based on type I string models. Our purpose is to study explicitly CP violation effects in models with non–universal gaugino masses and A–terms. Type I models can realize such initial conditions. These models contain nine–branes and three types of five–branes (5_a, $a = 1, 2, 3$). Here we assume that the gauge group $SU(3) \times U(1)_Y$ is on a 9–brane and the gauge group $SU(2)$ on the 5_1–brane like in Ref. [34, 43], in order to get non–universal gaugino masses between $SU(3)$ and $SU(2)$. We call these branes the $SU(3)$–brane and the $SU(2)$–brane, respectively.

Chiral matter fields correspond to open strings spanning between branes. Thus, they must be assigned accordingly to their quantum numbers. For example, the chiral field corresponding to the open string between the $SU(3)$ and $SU(2)$ branes has non–trivial representations under both $SU(3)$ and $SU(2)$, while the chiral field corresponding to the open string, which starts and ends on the $SU(3)$–brane, should be an $SU(2)$–singlet.

There is only one type of the open string that spans between the 9 and 5–branes, that we denote as the C^{95_1}. However, there are three types of open strings which start and end on the 9–brane, that is, the C^9_i sectors (i=1,2,3), corresponding to the i–th complex compact dimension among the three complex dimensions. If we assign the three families to the different C^9_i sectors we obtain non–universality in the right–handed sector. Notice that, in this model, we can not derive non–universality for the squark doublets, i.e. the left–handed sector. In particular, we assign the C^9_1 sector to the third family and the C^9_3 and C^9_2, to the first and second families, respectively.

Under the above assignment of the gauge multiplets and the matter fields, soft SUSY breaking terms are obtained, following the formulae in Ref. [21]. The gaugino masses are obtained

$$M_3 = M_1 = \sqrt{3} m_{3/2} \sin\theta e^{-i\alpha_S}, \tag{25}$$

$$M_2 = \sqrt{3} m_{3/2} \cos\theta \Theta_1 e^{-i\alpha_1}. \tag{26}$$

While the A–terms are obtained as

$$A_{C_1^9} = -\sqrt{3} m_{3/2} \sin\theta e^{-i\alpha_S} = -M_3, \tag{27}$$

for the coupling including C_1^9, i.e. the third family,

$$A_{C_2^9} = -\sqrt{3} m_{3/2} (\sin\theta e^{-i\alpha_S} + \cos\theta(\Theta_1 e^{-i\alpha_1} - \Theta_2 e^{-i\alpha_2})), \tag{28}$$

for the coupling including C_2^9, i.e. the second family and

$$A_{C_3^9} = -\sqrt{3} m_{3/2} (\sin\theta e^{-i\alpha_S} + \cos\theta(\Theta_1 e^{-i\alpha_1} - \Theta_3 e^{-i\alpha_3})), \tag{29}$$

for the coupling including C_3^9, i.e. the first family. Here $m_{3/2}$ is the gravitino mass, α_S and α_i are the CP phases of the F–terms of the dilaton field S and the three moduli fields T_i, and θ and Θ_i are goldstino angles, and we have the constraint, $\sum \Theta_i^2 = 1$.

Thus, if quark fields correspond to different C_i^9 sectors, we have non–universal A–terms. We obtain the following A–matrix for both of the up and down sectors,

$$A = \begin{pmatrix} A_{C_3^9} & A_{C_2^9} & A_{C_1^9} \\ A_{C_3^9} & A_{C_2^9} & A_{C_1^9} \\ A_{C_3^9} & A_{C_2^9} & A_{C_1^9} \end{pmatrix}. \tag{30}$$

The trilinear SUSY breaking matrix, $(Y^A)_{ij} = (Y)_{ij}(A)_{ij}$, itself is obtained

$$Y^A = \begin{pmatrix} & & \\ & Y_{ij} & \\ & & \end{pmatrix} \cdot \begin{pmatrix} A_{C_3^9} & 0 & 0 \\ 0 & A_{C_2^9} & 0 \\ 0 & 0 & A_{C_1^9} \end{pmatrix}, \tag{31}$$

in matrix notation.

In addition, soft scalar masses for quark doublets and the Higgs fields are obtained,

$$m^2_{C^{95_1}} = m^2_{3/2}(1 - \frac{3}{2}\cos^2\theta(1 - \Theta_1^2)). \quad (32)$$

The soft scalar masses for quark singlets are obtained as

$$m^2_{C_i^9} = m^2_{3/2}(1 - 3\cos^2\theta\Theta_i^2), \quad (33)$$

if it corresponds to the C_i^9 sector.

Now, below the string or SUSY breaking scale, this model is simply a MSSM with non–trivial soft–breaking terms from the point of view of flavor. Scalar mass matrices and tri–linear terms have completely new flavor structures, as opposed to the super–gravity inspired CMSSM or the SM, where the only connection between different generations is provided by the Yukawa matrices.

This model includes, in the quark sector, 7 different structures of flavor, M_Q^2, M_U^2, M_D^2, Y_d, Y_u, Y_d^A and Y_u^A. From these matrices, M_Q^2, the squark doublet mass matrix, is proportional to the identity matrix, and hence trivial, then we are left with 6 non–trivial flavor matrices. Notice that we have always the freedom to diagonalize the hermitian squark mass matrices (as we have done in the previous section, Eqs.(32,33)) and fix some general form for the Yukawa and tri–linear matrices. In this case, these four matrices are completely observable, unlike in the SM or CMSSM case.

At this point, to specify completely the model, we need not only the soft–breaking terms but also the complete Yukawa textures. The only available experimental information is the Cabbibo–Kobayashi–Maskawa (CKM) mixing matrix and the quark masses. Here, we choose our Yukawa texture following two simple assumptions : i) the CKM mixing matrix originates from the down Yukawa couplings (as done in the MM case) and ii) our Yukawa matrices are hermitian [44]. With these two assumptions we fix completely the Yukawa matrices,

$$Y_u = \frac{1}{v_2}\begin{pmatrix} m_u & 0 & 0 \\ 0 & m_c & 0 \\ 0 & 0 & m_t \end{pmatrix} \qquad Y_d = \frac{1}{v_1}\, K^\dagger \,.\begin{pmatrix} m_d & 0 & 0 \\ 0 & m_s & 0 \\ 0 & 0 & m_b \end{pmatrix}.\, K \quad (34)$$

with $v = v_1/(\cos\beta) = v_2/(\sin\beta) = \sqrt{2}M_W/g$, and K the CKM matrix. We take $\tan\beta = v_2/v_1 = 2$ in the following in all numerical examples. In this basis we can analyze the down tri–linear matrix,

$$Y_d^A(M_{St}) = \frac{1}{v_1}\ K^\dagger\ .\ M_d\ .\ K\ .\begin{pmatrix} A_{C_3^9} & 0 & 0 \\ 0 & A_{C_2^9} & 0 \\ 0 & 0 & A_{C_1^9} \end{pmatrix} \tag{35}$$

with $M_d = diag.(m_d, m_s, m_b)$.

Hence, together with the up tri–linear matrix we have our MSSM completely defined. The next step is simply to use the MSSM Renormalization Group Equations [27, 26] to obtain the whole spectrum and couplings at the low scale, M_W. The dominant effect in the tri–linear terms renormalization is due to the gluino mass which produces the well–known alignment among A–terms and gaugino phases. However, this renormalization is always proportional to the Yukawa couplings and not to the tri–linear terms, Eq.(7). This implies that, in the SCKM basis, the gluino effects will be diagonalized in excellent approximation, while due to the different flavor structure of the tri–linear terms large off–diagonal elements will remain with phases $\mathcal{O}(1)$ [37]. To see this more explicitly, we can roughly approximate the RGE effects as,

$$Y_d^A(M_W) = c_{\tilde{g}}\ m_{\tilde{g}}\ Y_d\ + c_A\ Y_d\ .\begin{pmatrix} A_{C_3^9} & 0 & 0 \\ 0 & A_{C_2^9} & 0 \\ 0 & 0 & A_{C_1^9} \end{pmatrix} \tag{36}$$

with $m_{\tilde{g}}$ the gluino mass and $c_{\tilde{g}}$, c_A coefficients order 1 (typically $c_{\tilde{g}} \simeq 5$ and $c_A \simeq 1$). We go to the SCKM basis after diagonalizing all the Yukawa matrices (that is, $K.Y_d.K^\dagger = M_d/v_1$). In this basis, we obtain the tri–linear couplings as,

$$v_1\ Y_d^A(M_W) = \Big(c_{\tilde{g}}\ m_{\tilde{g}}\ M_d\ + c_A\ M_d\ .\ K\ .\begin{pmatrix} A_{C_3^9} & 0 & 0 \\ 0 & A_{C_2^9} & 0 \\ 0 & 0 & A_{C_1^9} \end{pmatrix}\ .\ K^\dagger\Big) \tag{37}$$

From this equation we can get the L–R down squark mass matrix

$$m^2_{LR}{}^{(d)} = v_1\ Y_d^{A^*} - \mu e^{i\varphi_\mu} \tan\beta\, M_d \tag{38}$$

And finally using unitarity of K we obtain for the L–R Mass Insertions,

$$\begin{aligned}(\delta_{LR}^{(d)})_{ij} = \frac{1}{m^2_{\tilde{q}}}\ m_i\ \Big(\delta_{ij}\ (c_A A^*_{C_3^9}\ +\ c_{\tilde{g}}\ m^*_{\tilde{g}} -\ \mu e^{i\varphi_\mu} \tan\beta) + \\ K_{i2}\ K^*_{j2}\ c_A\ (A^*_{C_2^9} - A^*_{C_3^9}) + K_{i3}\ K^*_{j3}\ c_A\ (A^*_{C_1^9} - A^*_{C_3^9})\Big)\end{aligned} \tag{39}$$

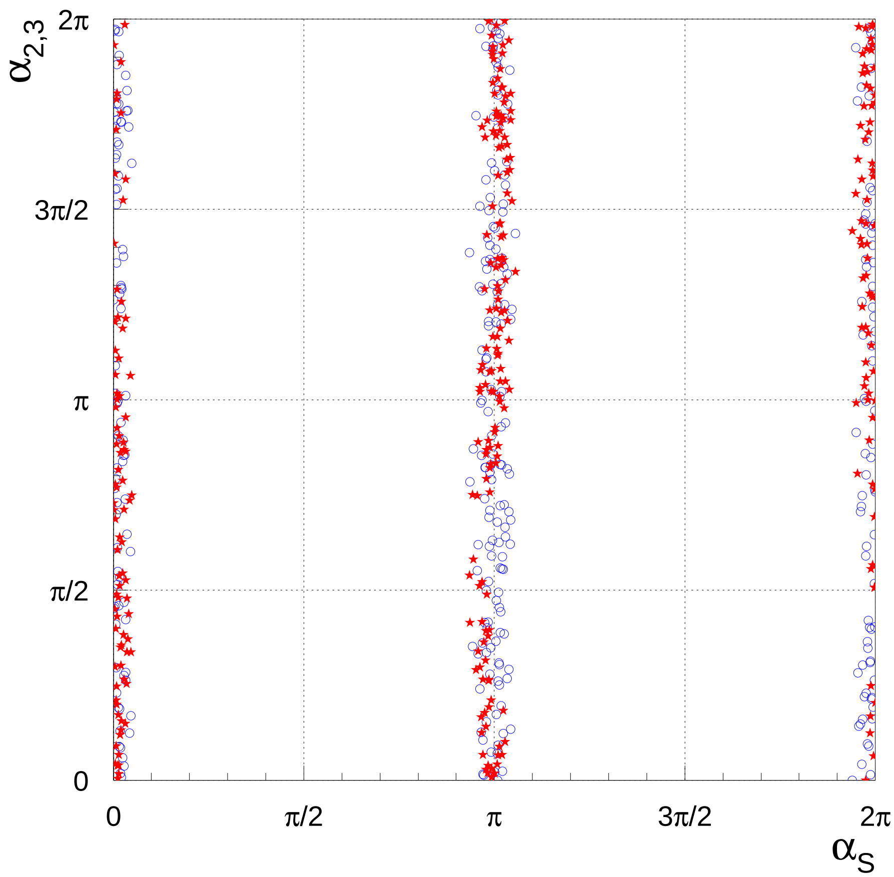

Figure 4: Allowed values for α_2–α_S (open blue circles) and α_3–α_S (red stars)

where $m_{\tilde{q}}^2$ is an average squark mass and m_i the quark mass. The same rotation must be applied to the L–L and R–R squark mass matrices,

$$M_{LL}^{(d)^2}(M_W) = K \ . \ M_Q^2(M_W) \ . \ K^\dagger$$
$$M_{RR}^{(d)^2}(M_W) = K \ . \ M_D^2(M_W) \ . \ K^\dagger \qquad (40)$$

However, the off–diagonal MI in these matrices are sufficiently small in this case thanks to the universal and dominant contribution from gluino to the squark mass matrices in the RGE. At this point, with the explicit expressions for $(\delta_{LR}^{(d)})_{ij}$, we can study the gluino mediated contributions to EDMs and ε'/ε. In this non–universal scenario, it is relatively easy to maintain the SUSY contributions to the EDM of the electron and the neutron below the experimental bounds while having large SUSY phases that contribute to ε'/ε. This is due to the fact the EDM are mainly controled by flavor–diagonal MI, while gluino contributions to ε'/ε are controled by $(\delta_{LR}^{(d)})_{12}$ and $(\delta_{LR}^{(d)})_{21}$. Here, we can have a very small phase for $(\delta_{LR}^{(d)})_{11}$ and $(\delta_{LR}^{(u)})_{11}$ and phases $\mathcal{O}(1)$ for the off–diagonal elements without any fine–tuning [42]. It is important to remember that the observable phase is always the relative phase between these mass insertions and the relevant gaugino mass involved. In Eq.(39) we can see that the diagonal elements tend to align with the gluino phase, hence to have a small EDM, it is enough to have the phases of the gauginos and the μ term approximately equal, $\alpha_S = \alpha_1 = -\varphi_\mu$. However α_2 and α_3 can still contribute to off–diagonal elements. In figure 4 we show the allowed values for α_S, α_2 and α_3 assuming $\alpha_1 = \varphi_\mu = 0$. We impose the EDM, ε_K and $b \to s\gamma$ bounds separately for gluino and chargino contributions together with the usual bounds on SUSY masses. We can see that, similarly to the CMSSM situation, φ_μ is constrained to be very close to the gluino and chargino phases (in the plot $\alpha_S \simeq 0, \pi$), but α_2 and α_3 are completely unconstrained.

Finally, in figure 5, we show the effects of these phases in the $(\delta_{LR}^{(d)})_{21}$ MI as a function of the gravitino mass. All the points in this plot satisfy all CP–conserving constraints besides EDM and ε_K constraints. We must remember that a value of $|\mathrm{Im}(\delta_{12}^d)_{LR}^2| \sim 10^{-5}$ gives a significant contribution to ε'/ε. In this plot, we can see a large percentage of points above or close to 1×10^{-5}. Hence, we can conclude that, in the presence of new flavor structures in the SUSY soft–breaking terms, it is not difficult to obtain sizeable SUSY

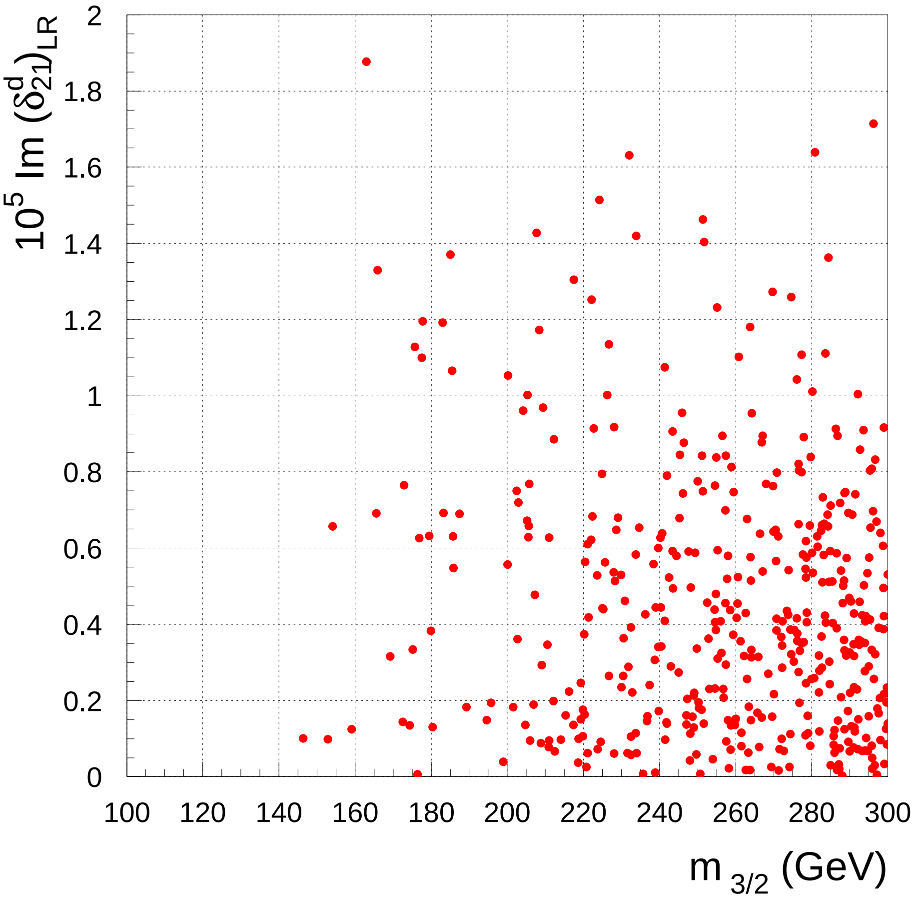

Figure 5: $(\delta_{LR}^{(d)})_{21}$ versus $m_{3/2}$ for experimentally allowed regions of the SUSY parameter space

contributions to CP violation observables and specially to ε'/ε [37, 42].[2]

5 Conclusions and Outlook

Here we summarize the main points of these lectures:

- There exist strong theoretical and "observational" reasons to go beyond the SM.
- The gauge hierarchy and coupling unification problems favor the presence of low–energy SUSY (either in its minimal version, CMSSM, or more naturally, in some less constrained realization).
- Flavor and CP problems constrain low–energy SUSY, but, at the same time, provide new tools to search for SUSY indirectly.
- In all generality, we expect new CP violating phases in the SUSY sector. However, these new phases are not going to produce sizeable effects as long as the SUSY model we consider does not exhibit a new flavor structure in addition to the SM Yukawa matrices.
- In the presence of a new flavor structure in SUSY, we showed that large contributions to CP violating observables are indeed possible.

In summary, given the fact that LEP searches for SUSY particles are close to their conclusion and that for Tevatron it may be rather challenging to find a SUSY evidence, we consider CP violation a potentially precious ground for SUSY searches before the advent of the "SUSY machine", LHC.

Acknowledgments

We thank D. Demir, T. Kobayashi, S. Khalil and H. Murayama as co–authors of some recent works reported in these lectures. We are grateful to S. Bertolini, L. Silvestrini and F.J. Botella for enlightening conversations. A.M. thanks the organizers for the stimulating settling in which the school

[2]With these L–R mass insertions alone, it is in general difficult to saturate ε_K [25]. However, in some special situations, it is still possible to have large contributions [34, 45]

took place. The work of A.M. was partly supported by the TMR project "Beyond the Standard Model" contract number ERBFMRX CT96 0090; O.V. acknowledges financial support from a Marie Curie EC grant (TMR-ERBFMBI CT98 3087).

References

[1] J. Ellis and D.V. Nanopoulos, *Phys. Lett.* **B 110** (1982) 44;
R. Barbieri and R. Gatto, *Phys. Lett.* **B 110** (1982) 211.

[2] For a phenomenologically oriented review, see:
P. Fayet and S. Ferrara, *Phys. Rep.* **32C** (1977) 249;
H.P. Nilles, *Phys. Rep.* **110C** (1984) 1;
H.E. Haber and G.L Kane, *Phys. Rep.* **117C** (1987) 1;
For spontaneously broken N=1 supergravity, see:
E. Cremmer, S. Ferrara, L. Girardello and A. Van Proeyen, *Nucl. Phys.* **B 212** (1983) 413 and references therein;
P. Nath, R. Arnowitt and A.H. Chamseddine, *Applied N=1 Supergravity* (World Scientific, Singapore, 1984);
A.G. Lahanas and D.V. Nanopoulos, *Phys. Rep.* **145C** (1987) 1.

[3] M. Dine, W. Fischler and M. Srednicki, *Nucl. Phys.* **B 189** (1981) 575;
S. Dimopoulos and S. Raby, *Nucl. Phys.* **B 192** (1981) 353;
M. Dine and W. Fischler, *Phys. Lett.* **B 110** (1982) 227;
M. Dine and M. Srednicki, *Nucl. Phys.* **B 202** (1982) 238;
M. Dine and W. Fischler, *Nucl. Phys.* **B 204** (1982) 346;
L. Alvarez–Gaumé, M. Claudson and M. Wise, *Nucl. Phys.* **B 207** (1982) 96;
C. Nappi and B. Ovrut, *Phys. Lett.* **B 113** (1982) 175;
S. Dimopoulos and S. Raby, *Nucl. Phys.* **B 219** (1983) 479.

[4] A. Nelson and M. Dine, *Phys. Rev.* **D 48** (1993) 1277, hep-ph/9303230;
M. Dine, A. E. Nelson and Y. Shirman, *Phys. Rev.* **D 51** (1995) 1362, hep-ph/9408384;
M. Dine, A. Nelson, Y. Nir and Y. Shirman, *Phys. Rev.* **D 53** (1996) 2658, hep-ph/9507378.

[5] E. Poppitz and S. Trivedi, *Phys. Rev.* **D 55** (1997) 5508, hep-ph/9609529;
N. Arkani–Hamed, J. March–Russel and H. Murayama, *Nucl. Phys.* **B 509** (1998) 3, hep-ph/9701286;
H. Murayama, *Phys. Rev. Lett.* **79** (1997) 18, hep-ph/9705271;
S. Dimopoulos, G. Dvali, G. Giudice and R. Rattazzi, *Nucl. Phys.* **B 510** (1998) 12, hep-ph/9705307;
S. Dimopoulos, G. Dvali and R. Rattazzi, *Phys. Lett.* **B 413** (1997) 336, hep-ph/9707537 ;
M. Luty, *Phys. Lett.* **B 414** (1997) 71, hep-ph/9706554;
T. Hotta, K.-I. Izawa and T. Yanagida, *Phys. Rev.* **D 55** (1997) 415, hep-ph/9606203;
N. Haba, N. Maru and T. Matsuoka, *Nucl. Phys.* **B 497** (1997) 31, hep-ph/9612468;
L. Randall, *Nucl. Phys.* **B 495** (1997) 37, hep-ph/9612426;
Y. Shadmi, *Phys. Lett.* **B 405** (1997) 99, hep-ph/9703312;
N. Haba, N. Maru and T. Matsuoka, *Phys. Rev.* **D 56** (1997) 4207, hep-ph/9703250;
C. Csaki, L. Randall and W. Skiba, *Phys. Rev.* **D 57** (1998) 383, hep-ph/9707386;
Y. Shirman, *Phys. Lett.* **B 417** (1998) 281, hep-ph/9709383;
G.F. Giudice and R. Rattazzi, hep-ph/9801271.

[6] M. Dugan, B. Grinstein and L. Hall, *Nucl. Phys.* **B 255** (1985) 413.

[7] S. Dimopoulos and S. Thomas, *Nucl. Phys.* **B 465** (1996) 23, hep-ph/9510220.

[8] W. Buchmuller and D. Wyler, *Phys. Lett.* **B 121** (1983) 321;
J. Polchinski and M. Wise, *Phys. Lett.* **B 125** (1983) 393;
W. Fischler, S. Paban and S. Thomas, *Phys. Lett.* **B 289** (1992) 373, hep-ph/9205233.

[9] J. Ellis and R. Flores, *Phys. Lett.* **B 377** (1996) 83, hep-ph/9602211.

[10] M. Dine, A. Nelson and Y. Shirman, in ref. [4];
M. Dine, Y. Nir and Y. Shirman, *Phys. Rev.* **D 55** (1997) 1501, hep-ph/9607397.

[11] M. J. Duncan and J. Trampetic, *Phys. Lett.* **B 134** (1984) 439;
E. Franco and M. Mangano, *Phys. Lett.* **B 135** (1984) 445;
J.M. Gerard, W. Grimus, A. Raychaudhuri and G. Zoupanos, *Phys. Lett.* **B 140** (1984) 349;
J. M. Gerard, W. Grimus, A. Masiero, D. V. Nanopoulos and A. Raychaudhuri, *Phys. Lett.* **B 141** (1984) 79; *Nucl. Phys.* **B 253** (1985) 93;
P. Langacker and R. Sathiapalan, *Phys. Lett.* **B 144** (1984) 401;
M. Dugan, B. Grinstein and L. Hall, in ref. [6].

[12] A. Brignole, F. Feruglio and F. Zwirner, *Zeit. für Physik* **C 71** (1996) 679, hep-ph/9601293.

[13] M. Misiak, S. Pokorski and J. Rosiek, in *Heavy Flavours II*, eds. A. Buras and M. Lindner (World Scientific), hep-ph/9703442.

[14] G.C. Branco, G.C. Cho, Y. Kizukuri and N. Oshimo, *Nucl. Phys.* **B 449** (1995) 483;
G.C. Branco, G.C. Cho, Y. Kizukuri and N. Oshimo, *Phys. Lett.* **B 337** (1994) 316, hep-ph/9408229.

[15] T. Ibrahim and P. Nath, *Phys. Rev.* **D 58** (1998) 111301, hep-ph/9807501;
M. Brhlik, G.J. Good and G.L. Kane, *Phys. Rev.* **D 59** (1999) 115004, hep-ph/9810457;
A. Bartl, T. Gajdosik, W. Porod, P. Stockinger and H. Stremnitzer, *Phys. Rev.* **D 60** (1999) 073003, hep-ph/9903402.

[16] S.A. Abel and J.M. Frere, *Phys. Rev.* **D 55** (1997) 1623, hep-ph/9608251;
S. Khalil, T. Kobayashi and A. Masiero, *Phys. Rev.* **D 60** (1999) 075003, hep-ph/9903544;
S. Khalil and T. Kobayashi, *Phys. Lett.* **B 460** (1999) 341, hep-ph/9906374.

[17] S. Dimopoulos and G.F. Giudice, *Phys. Lett.* **B 357** (1995) 573, hep-ph/9507282;
A. Cohen, D.B. Kaplan and A.E. Nelson, *Phys. Lett.* **B 388** (1996) 599, hep-ph/9607394;

A. Pomarol and D. Tommasini, *Nucl. Phys.* **B 466** (1996) 3, hep-ph/9507462.

[18] Y. Grossman, Y. Nir and R. Rattazzi, in *Heavy Flavours II*, eds. A. Buras and M. Lindner (World Scientific), hep-ph/9701231.

[19] D. Demir, A. Masiero and O. Vives, SISSA report n. SISSA/134/99/EP, hep-ph/9911337.

[20] D. Demir, A. Masiero and O. Vives, *Phys. Rev.* **D 61** (2000) 075009, hep-ph/9909325.

[21] L.E. Ibanez, C. Munoz and S. Rigolin, *Nucl. Phys.* **B 553** (1999) 43, hep-ph/9812397.

[22] M. Brhlik, L. Everett, G.L. Kane and J. Lykken, *Phys. Rev. Lett.* **83** (1999) 2124, hep-ph/9905215;
M. Brhlik, L. Everett, G.L. Kane and J. Lykken, Fermilab preprint no. FERMILAB-PUB-99-230-T, August 1999, hep-ph/9908326;
T. Ibrahim and P. Nath, Santa Barbara U. preprint no. NSF-ITP-99-129, October 1999, hep-ph/9910553.

[23] D. Demir, A. Masiero and O. Vives, *Phys. Rev. Lett.* **82** (1999) 2447, Err. *ibid.* **83** (1999) 2093, hep-ph/9812337.

[24] S. Baek and P. Ko, *Phys. Rev. Lett.* **83** (1999) 488, hep-ph/9812229;
S. Baek and P. Ko, Korea Inst. Sci Report No. KAIST-TH-99-1, 1999, hep-ph/9904283.

[25] F. Gabbiani, E. Gabrielli, A. Masiero and L. Silvestrini, *Nucl. Phys.* **B 477** (1996) 321, hep-ph/9604387.

[26] S. Bertolini, F. Borzumati, A. Masiero and G. Ridolfi, *Nucl. Phys.* **B 353** (1991) 591.

[27] N.K. Falck, *Zeit. für Physik* **C 30** (1986) 247.

[28] L.J. Hall, V.A. Kostelecky and S. Raby, *Nucl. Phys.* **B 267** (1986) 415.

[29] H.E. Haber and G.L Kane in ref. [2].

[30] P. Cho, M. Misiak and D. Wyler, *Phys. Rev.* **D 54** (1996) 3329, hep-ph/9601360.

[31] F.M. Borzumati, *Zeit. für Physik* **C 63** (1994) 291, hep-ph/9310212;
S. Bertolini and F. Vissani, *Zeit. für Physik* **C 67** (1995) 513, hep-ph/9403397;
T. Goto, Y.Y. Keum, T. Nihei, Y. Okada and Y. Shimizu, *Phys. Lett.* **B 460** (1999) 333, hep-ph/9812369.

[32] A.L. Kagan and M. Neubert, *Eur. Phys. J.* **C7** (1999) 5, hep-ph/9805303;
A.L. Kagan and M. Neubert, *Phys. Rev.* **D 58** (1998) 094012, hep-ph/9803368.

[33] M. Ciuchini *et al. J. High Energy Phys.* **10** (98) 008, hep-ph/9808328;
R. Contino and I. Scimemi, *Eur. Phys. J.* **C10** (1999) 347, hep-ph/9809437.

[34] M. Brhlik, L. Everett, G.L. Kane, S.F. King and O. Lebedev, Virginia Pol. Inst. Report no. VPI/IPPAP/99/08, Sept. 1999, hep-ph/9909480;

[35] R. Barbieri, R. Contino and A. Strumia, Pisa U. Report no. IFUP-TH-45-99, Aug. 1999, hep-ph/9908255.
A. L. Kagan and M. Neubert, *Phys. Rev. Lett.* **83** (1999) 4429, hep-ph/9908404.
K. S. Babu, B. Dutta and R. N. Mohapatra, Oklahoma State U. preprint no. OSU-HEP-99-03, hep-ph/9905464.

[36] E. Gabrielli and G. F. Giudice, *Nucl. Phys.* **B 433** (1995) 3, hep-lat/9407029.

[37] A. Masiero and H. Murayama, *Phys. Rev. Lett.* **83** (1999) 907, hep-ph/9903363.

[38] See, *e.g.*, Y. Kawamura, H. Murayama, and M. Yamaguchi, *Phys. Rev.* **D 51** (1995) 1337, hep-ph/9406245;
A. Brignole, L.E. Ibañez, and C. Muñoz, in *Perspectives on Supersymmetry*, *ed.* G. Kane, Singapore, World Scientific, 1998, hep-ph/9707209.

[39] KTeV Collaboration, A. Alavi-Harati *et al.*, *Phys. Rev. Lett.* **83** (1999) 22, hep-ex/9905060.

[40] NA31 Collaboration (G.D. Barr *et al.*), *Phys. Lett.* **B 317** (1993) 233.

[41] S.A. Abel and J.M. Frere in ref. [16].

[42] S. Khalil, T. Kobayashi and O. Vives, SISSA report number SISSA/23/20000/EP, hep-ph/0003086.

[43] T. Ibrahim and P. Nath in ref. [22]

[44] P. Ramond, R. G. Roberts and G. G. Ross, *Nucl. Phys.* **B 406** (1993) 19, hep-ph/9303320.

[45] G. D'Ambrosio, G. Isidori and G. Martinelli, INFN preprint no. INFN-NA-99-39, hep-ph/9911522.